arXiv:1304.2767v2 [hep-th] 11 Apr 2013

The Higgs mass from a String-Theoretic Perspective

Arthur Hebecker[1], Alexander K. Knochel[1,2] and Timo Weigand[1]

[1]*Institut für Theoretische Physik, Universität Heidelberg, Philosophenweg 19,*
D-69120 Heidelberg, Germany

[2]*Institut für Theoretische Teilchenphysik und Kosmologie, RWTH Aachen,*
D-52056 Aachen, Germany

(a.hebecker , t.weigand@thphys.uni-heidelberg.de ;
knochel@physik.rwth-aachen.de)

Abstract

The Higgs quartic coupling λ has now been indirectly measured at the electroweak scale. Assuming no new low-scale physics, its running is known and, together with gauge and Yukawa couplings, it is a crucial new piece of information constraining UV completions of the Standard Model. In particular, supersymmetry broken at an intermediate or high energy scale with $\tan\beta = 1$ (i.e. $\lambda = 0$) is consistent with present data and has an independent theoretical appeal. We analyze the possible string-theoretic motivations for $\tan\beta = 1$ (including both the shift-symmetry and the more economical variant of a Z_2 symmetry) in a Higgs sector realized on either 6- or 7-branes. We identify specific geometries where $\lambda \simeq 0$ may arise naturally and specify the geometrical problems which need to be solved to determine its precise value in the generic case. We then analyze the radiative corrections to λ. Finally we show that, in contrast to naive expectations, $\lambda < 0$ at the SUSY breaking scale is also possible. Specifically, string theory may produce an MSSM plus chiral singlet at a very high scale, which immediately breaks to a non-SUSY Standard Model with $\lambda < 0$. This classically unstable theory then becomes metastable through running towards the IR.

1   Introduction

It is now established that the recently discovered resonance described in [1, 2] is at least very similar [77] to the Standard Model Higgs boson. Together with constraints from electroweak precision data, flavor physics (see e.g. [3, 4] and refs. therein) and quickly evolving direct-search limits from LHC, this supports the conservative hypothesis that the Standard Model remains the correct effective field theory above the TeV scale. To the best of our knowledge, this implies that the electroweak scale is fine-tuned, which may indeed be acceptable in view of the huge landscape of string vacua (see e.g. [5]).

Thus, we focus on string theory as *the* high-scale theory and, more specifically, on those solutions of the string equations of motion which can be understood as compactifications of 10d supergravity (F-theory models, heterotic Calabi-Yaus, or type II Calabi-Yau orientifolds with branes).[1] For reasons of stability, we expect supersymmetry at the compactification scale m_C. If we accept fine tuning in the Higgs sector, SUSY can then be broken anywhere below that scale, $m_Z \ll m_S \leq m_C$. For the purpose of this paper, we mostly focus on the 'canonical' minimal particle content at m_S: a high-scale MSSM. We furthermore assume that all superpartner masses are of the same order of magnitude (see however [6]), such that m_S is reasonably well-defined.

Now, given the measured Higgs mass, the electroweak-scale quartic coupling λ of the Higgs field and its running at higher energies are in principle known. In particular, λ runs to zero at some high scale μ_λ [7–12] (see [13] for state-of-the-art SM beta functions). The standard MSSM formula

$$\lambda(m_S) = \frac{g_2^2(m_S) + g_1^2(m_S)}{8} \cos^2 2\beta \qquad (1)$$

expresses the quartic coupling in terms of the MSSM β-angle and the electroweak gauge couplings g_1 and g_2 [14]. One sees that $\lambda(m_S) \geq 0$, implying $m_S \leq \mu_\lambda$. Furthermore, if an argument for a particular value of $\tan\beta$ can be made, Eq. (1) together with the running of λ can be used to predict m_S.[2] This was exploited in [15], with focus on $\tan\beta \gg 1$ at the GUT scale (now ruled out by data). If, by contrast, a convincing theory reason for the opposite limit, $\tan\beta = 1$ and thus $\cos^2 2\beta = 0$, can be provided, high-scale SUSY is consistent with a 126 GeV Higgs and the SUSY breaking scale m_S coincides with the scale μ_λ at which the quartic coupling λ vanishes. In this scenario, introduced in [16], the prediction $m_S \sim \mu_\lambda \sim 10^9$ GeV (at present still with very large errors) can be made. More specifically, in [16] a string-motivated [17, 18] shift symmetry was suggested as the origin of $\tan\beta = 1$.[3] As observed in [22] (and further analyzed phenomenologically in [23]), the weaker assumption of a Z_2 symmetry is in fact sufficient to ensure $\tan\beta = 1$.

---

[1] Of course, entirely stringy constructions without any higher-dimensional interpretation and with supersymmetry broken at the string scale are also conceivable. The pragmatic reason for not considering them here is our lack of quantitative understanding. A possible physics reason is their short expected life-time: More or less by definition, there is no small parameter controlling tunneling rates in the corresponding part of the landscape.

[2] For a detailed numerical study and plots illustrating this relation see e.g. [9].

[3] Non SUSY-variants of a shift symmetry in the Higgs sector appeared much earlier in [19] as well as very recently [20] in a related context. An alternative origin of $\tan\beta = 1$ was suggested in [21].

There is, however, an interesting alternative to Eq. (1) in non-minimal SUSY models which allows $\lambda < 0$ at the soft scale. Note that this is independent of the details of the stringy or supersymmetric UV completion. We will introduce this scenario in section 4.2 and argue that an energy window *without* even a metastable vacuum at small or zero Higgs vev might exist. Rather than extending the SM by a stabilizing sector below μ_λ [11, 24], we simply accept that λ runs negative and only ensure the existence of some UV completion at much higher energies.

In the present paper, we attempt to flesh out the possible stringy symmetries in the Higgs sector (either shift or Z_2) which were suggested in [16, 22] as the origin of $\tan\beta = 1$. While the shift symmetry is well-understood in heterotic orbifold models [17, 18], its possible generalization to heterotic Calabi-Yaus appears to be complicated and we postpone this issue. Instead, we focus on 6- and 7-brane models in type IIA and IIB Calabi-Yau orientifolds and in F-theory. We also compare shift symmetry violating effects arising from the inequality $m_C > m_S$ and from loop corrections at m_S.

Before giving an outline of the paper, we now pause to discuss critically possible motivations for high-scale SUSY with $\tan\beta = 1$. The simplest motivation, already emphasized above, is that it may relate a symmetry feature of the high-scale theory (Z_2 or shift symmetry) to the last observable of the Standard Model (Higgs mass or equivalently quartic coupling λ). As such, scenarios with $\tan\beta = 1$ are certainly worthy of investigation.

Furthermore, it is conceivable that a better understanding of 'how to find a Standard Model in the landscape' will provide us with a first-principle reason for a symmetry enforcing $\tan\beta = 1$. As will become more apparent below, this is far from obvious. In fact, at the moment it appears that the required symmetries are only available in a certain class of models and that (as usual) many models *without* this feature exist.

Finally, there might be a 'landscape bias' towards a high SUSY breaking scale. Given a certain low-scale value of λ with a corresponding 'vacuum stability scale' $\mu_\lambda < M_{\mathrm{P}}$, this implies a bias for m_S to be as close to μ_λ as possible[4]. In such a situation, models ensuring $\tan\beta = 1$ by a symmetry are very strongly favored. While there are obviously many uncertainties in this argument (e.g. the preferred SUSY breaking scale as well as the whole concept of a measure in the landscape), we still believe that it is an interesting extra reason to consider the present class of models.

The paper is organized as follows: In the remainder of the introduction, we briefly recall the role of shift and exchange symmetry in determining the UV boundary conditions of the Higgs sector masses and couplings.

In Section 2, we discuss possible stringy realizations of shift/exchange symmetric Higgs sectors in smooth Calabi-Yau compactifications. In compactifications with D-branes, the Higgs field can either propagate along the entire brane as a so-called bulk Higgs, or localize at the the intersection of two D-branes. We argue that a bulk Higgs on D6 branes or its type IIB/F-Theory dual might be the most natural way to obtain shift

---

[4]Note that there is an unconventional alternative if one is willing to accept $\lambda < 0$ at the matching scale and possible metastability of the SM vacuum. This possibility is discussed in section 4.2.

symmetric scenarios in type II models. Upon deforming parallel branes relative to each other, bulk degrees of freedom become localized matter. We take a closer look at the fate of shift/exchange symmetry during the transition from D7 bulk fields to matter localized on the intersection curve of two D7-branes, and identify the point in moduli space where the transition between the two regimes occurs. We consider simplified settings in $D = 5, 6$ with modulus dominated or Scherck-Schwarz supersymmetry breaking and then draw some conclusions for Calabi-Yau compactifications with fluxes.

In Section 3, we discuss the Higgs quartic coupling in models where extended SUSY effects can play a significant role. In particular, we are interested in the recently proposed exchange symmetric models [22, 23] on D6 brane networks exploiting the possibility of relative extended $\mathcal{N} = 2$ supersymmetry preserved by pairs of D6 brane stacks. We give field-theoretic arguments that there is some tension between the need to generate a sizable $B\mu$ term via branes at angles or fluxes, and the desire to have an MSSM-like D-term potential. We discuss variations of the branes-at-angles scenario in the flat space limit and propose a possible solution.

Section 4 addresses both tree level and loop corrections to the quartic Higgs coupling and the intriguing possibility of negative λ at the soft scale. In Section 4.1, our previous analysis of radiative shift symmetry violation is updated and augmented by a detailed treatment of 1-loop threshold effects. We discuss the relative importance of the various contributions and their impact on the predicted SUSY breaking scale. In Section 4.2, we argue that some high-scale UV completions with approximately flat directions (e.g. approximately shift- or exchange symmetric models) can yield negative effective quartic couplings and nevertheless remain calculable perturbatively due to the local flatness of the potential at small field values. We show how this relaxes the upper bounds on the SUSY breaking scale, which might be useful for certain classes of string models (see e.g. [25]).

1.1 The role of shift symmetry and exchange symmetry

Let us review the role of the shift and the exchange symmetries in the UV boundary conditions of the Higgs quartic couplings before we discuss superstring/F-Theory realizations.

Shift symmetry: We consider an MSSM-like Higgs sector with two chiral Higgs doublets H_u and H_d carrying opposite hypercharge. We demand that the Kähler potential exhibit a shift symmetry

$$H_u \longrightarrow H_u + c, \quad H_d \longrightarrow H_d - c^\dagger. \tag{2}$$

The consequence of this symmetry is that the Kähler potential can only depend on the linear combination, $K = K(H_u + H_d^\dagger)$. The lowest-dimension allowed operator in the Higgs fields is

$$K = f(X_i, \overline{X}_i)|H_u + H_d^\dagger|^2, \tag{3}$$

where we have included a generic dependence on some moduli X_i which will later acquire an F-term vev. Furthermore, the dominant source of shift symmetry violation should come from the couplings, i.e. the quadratic part of the Higgs lagrangian will be affected by this violation at the 1-loop level (see section 4.1 for a discussion of these effects). What is remarkable about this Kähler potential is that it contains the usual kinetic terms $|H_u|^2$, $|H_d|^2$ as well as a "holomorphic" part $H_u H_d + H_d^\dagger H_u^\dagger$. The latter by itself would only contribute a total derivative to the action, but this changes in the supergravity setting, and specifically due to the nontrivial dependence of f on the moduli.

In modulus dominated SUSY breaking, some of the moduli X_i will acquire F-term vevs F^{Xi}, which in turn can generate soft terms as well as supersymmetric masses wherever they occur in the Kähler potential and superpotential. In absence of a $\mu H_u H_d$ term in the superpotential, the Higgs mass matrix can be read off the Kähler potential in Eq. (3). Without loss of generality, we take $f = 1$ in the vacuum. Upon supersymmetry breaking, soft masses $m_{H_u}^2, m_{H_d}^2$, the $B\mu$ term and the effective μ term are generated. For simplicity we now assume one modulus X to dominate. The resulting Higgs mass matrix (with $B\mu \equiv m_3^2$) at the scale m_S, defined through

$$\mathcal{L} \supset -m_1^2|H_u|^2 - m_2^2|H_d|^2 - m_3^2(H_u H_d + H_d^\dagger H_u^\dagger) \,, \tag{4}$$

is then given by (cf. e.g. [17, 18])

$$m_1^2 = m_2^2 = m_3^2 = |\mu|^2 + m_{3/2}^2 - F^X \overline{F}^X (\ln f)_{X\overline{X}} \,, \tag{5}$$

where

$$|\mu|^2 = \left| m_{3/2} - \overline{F}^X f_{\overline{X}} \right|^2 \,, \qquad \overline{F}^X = e^{K/2} K^{\overline{X}X} D_S W \qquad \text{and} \qquad m_{3/2} = e^{K/2} W \,. \tag{6}$$

As was discussed in [16], this Higgs mass matrix has a massless eigenstate

$$H_0 = \frac{1}{\sqrt{2}}(H_u - H_d^\dagger) \,. \tag{7}$$

It provides the SM Higgs doublet, while the orthogonal combination becomes heavy at the soft scale and does not contribute to EWSB. This situation corresponds to the decoupling limit with a mixing angle $\tan\beta = 1$. Since we get a zero eigenvalue, it appears as though we get a naturally low (compared to the soft scale) electroweak scale and have avoided the hierarchy problem. Unfortunately, this changes once one goes beyond tree level. After integrating out heavy degrees of freedom coupling to the Higgs sector like the top squarks, the corresponding decoupling contributions, though loop suppressed, give us large shifts in the Higgs mass parameters which force us to finetune the electroweak scale. The fact that these contributions are loop suppressed with respect to the treelevel mass parameters ensures that $\tan\beta \approx 1$ still holds to good accuracy after loop corrections are taken into account.

One might worry that the shift-symmetric scenario is too restrictive to tune the Higgs light after threshold corrections at m_S are included. However, since the shift-symmetry violating contributions themselves are enhanced relative to the quadratic threshold corrections by a $\log m_S/m_C$, one can tune the Higgs light within the leading log approximation by dialing the soft terms without having to give up the shift-symmetric scenario [16].

Exchange symmetry: One can ask whether there are weaker conditions than a shift-symmetric Higgs sector which also yield the desired structures. In [22], the authors find that there is a simpler way to enforce $\tan\beta = 1$ by symmetry if one does not insist on a naturally small electroweak scale at tree level. Since one needs to finetune anyways after taking into account quadratic loop corrections, we do only a factor of $\approx 6y_t^2/16\pi^2$ worse if we give up this requirement altogether, as was also pointed out in [23]. The idea is then as follows: if the Higgs potential exhibits an exchange symmetry (again, necessarily only at tree level) between H_u and $H_d^\dagger$, this ensures that the diagonal entries in the Higgs mass matrix satisfy $m_{H_u}^2 = m_{H_d}^2 = m^2$, and thus after canonical normalization, it becomes

$$M(m_S) = \left[\begin{array}{cc} |\mu|^2 + m^2 & B\mu \\ B\mu & |\mu|^2 + m^2 \end{array} \right]. \tag{8}$$

The electroweak scale is tuned light when

$$B\mu \approx |\mu|^2 + m^2 + \text{quadratic thresholds}. \tag{9}$$

Thus, the requirement of having a light electroweak scale already implies that $m_3 \approx m_1 = m_2$, and we recover the same Higgs mass matrix with universal entries as before and thus $\tan\beta \approx 1$ along with a massless doublet. If the relation $m_{H_u}^2 = m_{H_d}^2$ receives small corrections, the resulting deviation from $\tan\beta = 1$ and thus $\cos 2\beta = 0$ is

$$|\cos 2\beta| = \frac{1}{2} \frac{|m_{H_u}^2 - m_{H_d}^2|}{|B\mu|} \tag{10}$$

at leading order in $\Delta m^2/\mu^2$. In a sense, the shift-symmetric scenario is a special case of the exchange symmetric one, and both symmetries are equally broken by loop effects from terms such as the large third generation Yukawa coupling, which for $\tan\beta = 1$ is only present for H_u, but not H_d. Our treatment of radiative corrections in section 4.1, which is an extension of the analysis given in [16], is thus also valid for this case. This is also in agreement with a recent analysis in [23]. A particularly simple version of the exchange symmetric scenario is realized if the soft masses m_{H_u}, m_{H_d} in the Higgs sector vanish altogether or are strongly suppressed relative to $|\mu|^2$. In [22], such a scenario is proposed using brane angles or open string fluxes to generate a $B\mu$ term. Another possible realization of this idea, which we will not consider further in this work, assumes that supersymmetry breaking is communicated to the Higgs sector via a renormalizable interaction $\mathcal{W} \sim S H_u H_d$ with a "PQ-spurion". If $\langle S \rangle = s + \theta^2 F^S$, the resulting $\mu/B\mu$ terms, which yield exchange symmetric normalized masses, might dominate over exchange-violating soft masses originating from higher-dimensional operators. The Higgs mass matrix then has the desired approximate exchange-symmetric form, and the electroweak scale can be tuned light by varying μ_{eff} versus $B\mu$. However, the scalar F-term will also contribute to the quartic Higgs interaction at tree level if S is not decoupled in an approximately supersymmetric fashion above the scale m_S, and this can spoil the Higgs mass predictions from $\tan\beta = 1$ (see section 3). Also, if this mechanism yields large hierarchies between the heavy Higgs doublet and the sfermions or gauginos, the threshold corrections to λ may be the dominant effect.

2   Type IIB/F theory with 7-branes

The earliest and most explicit models with a shift-symmetric Higgs doublet pair are heterotic orbifolds [17, 26]. As explained in some detail in [16, 27, 28], the shift symmetry emerges since the Higgs is a Wilson line [30] (see [29] for a recent discussion of twisted vs. untwisted Higgs sectors in heterotic orbifolds). Here, we are interested in generic, smooth Calabi-Yau compactifications. This generalization is problematic [16] since, in going from orbifold to Calabi-Yau, the Higgs becomes one of the bundle moduli and the corresponding moduli space is less well understood.

2.1   Shift symmetry: Wilson-lines vs. brane-deformations

We therefore turn to brane models, for concreteness type IIB with D7-branes (see e.g. [31] for constructions on orbifolds and smooth Calabi-Yau manifolds), where the Higgs once again has the chance of being simply a Wilson line. A strongly simplified version of the corresponding Kähler potential [32] is

$$K = -\ln[-i(S - \overline{S}) - L(z,\overline{z})\zeta\overline{\zeta}] - 3\ln[T + \overline{T} - C(z,\overline{z},\zeta,\overline{\zeta})a\overline{a}] - K_{cs}(z,\overline{z}) \cdots . \quad (11)$$

Here S is the axio-dilaton and z, ζ, T, a are the complex structure, D7-brane, Kähler and Wilson line moduli respectively. Merely for notational simplicity, we suppress indices pretending that there were just one modulus of each kind.

Since this Kähler potential has only been obtained in a fluctuation-expansion around a given brane configuration, we a priori do not know how the leading terms $\zeta\overline{\zeta}$ and $a\overline{a}$ generalize to all-orders Kähler metrics on the brane-deformation moduli space and the Wilson-line moduli-space respectively.

Naively, one might hope that $a\overline{a}$ actually arises from a term $\sim (a + \overline{a})^2$ or, more generally, from a *shift-symmetric* Wilson line Kähler potential

$$k_w(z,\overline{z},\zeta,\overline{\zeta},a+\overline{a}) \;=\; C(z,\overline{z},\zeta,\overline{\zeta})\, a\overline{a} + \cdots . \quad (12)$$

However, a closer look reveals that things can not be so simple: Our complex variable a combines two independent Wilson lines corresponding to its real and imaginary part. Respecting a shift-symmetry in both of them would imply that k_w should be independent of both $(a - \overline{a})$ and $(a + \overline{a})$, which is clearly impossible. Indeed, upon closer inspection of the dimensional reduction carried out in [32] one sees that the Chern-Simons term contains (very schematically) a piece

$$\int_{D7} C_4 \wedge dA_1 \wedge dA_1 = \int_{D7} A_1 \wedge dC_4 \wedge dA_1 , \quad (13)$$

which induces kinetic mixing between D7 Wilson lines and C_4 scalars, with a prefactor depending on the Wilson-line. This destroys the shift symmetry (at least generically).

Of course, even in the absence of a shift symmetry, an exchange symmetry of the type suggested in [22, 23] might still apply. Thus, we by no means rule out D7-brane

Wilson-line scalars as candidates for a Higgs sector with $\tan\beta = 1$. However, we will not develop this line of thinking in the present paper.

Next consider the brane deformation moduli $\zeta, \overline{\zeta}$ where, as we will argue, a shift symmetry does indeed exist (cf. [16,33,34]). To make our point, we consider the type IIA mirror version of (11). The analogue of the first term, again strongly simplified, is [35]

$$K = -\ln[-i(S - \overline{S}) - Q(t, \overline{t})u\overline{u}] + \cdots , \tag{14}$$

where t and u are the Kähler and D6-brane moduli respectively. To avoid losing focus, we do not discuss the type IIA analogue of S in any detail. We only note that the type IIA axio-dilaton contains both g_s and the volume and also has an intimate relation to the complex structure through the C_3 scalar (its imaginary part) associated with the holomorphic 3-form [36]. This will not be important for us and we do not go into the details of how S is expressed in terms of IIA compactification data. The claim that such a superfield exists and the Kähler potential can be written as in (14) can be viewed as the assertion that mirror symmetry continues to hold in the context of $\mathcal{N} = 1$ orientifolds.

The main point for us is that the complex field

$$u = \Phi + i(A - M\Phi) \tag{15}$$

combines one real brane-deformation Φ and one Wilson line modulus A (M vanishes for vanishing B_2 background and will not be relevant for us). In particular, $\text{Re}(u)$ is independent of A. Furthermore, as can be seen from [35], the dimensionally reduced 4d action involves $A \sim (u - \overline{u})$ only through its derivatives. This crucially depends on the structure of the relevant Chern-Simons term,

$$\int_{D6} C_3 \wedge dA_1 \wedge dA_1 = \int_{D6} A_1 \wedge dC_3 \wedge dA_1 . \tag{16}$$

In contrast to the D7-brane case discussed earlier, this term does not give rise to any (potentially A-dependent) kinetic mixing between the Wilson line scalar and a C_3 scalar. The point is that, for such a mixing term, both $A_1 s$ should be integrated along 1-cycles of the D6. For a non-zero result, one would need to dimensionally reduce dC_3 to a 4d 3-form using a harmonic Calabi-Yau 1-form, to be then pulled back to the brane. Since Calabi-Yaus have no 1-cycles, no such kinetic mixing terms arise.

Thus, we can generalize (14) by writing

$$K = -\ln[-i(S - \overline{S}) - k_{D6}(t, \overline{t}, u + \overline{u})] + \cdots , \tag{17}$$

or, returning to the IIB side,

$$K = -\ln[-i(S - \overline{S}) - k_{D7}(z, \overline{z}, \zeta + \overline{\zeta})] + \cdots . \tag{18}$$

Note that we need to be at large complex structure on the type IIB side in order to trust the classical Kähler potential found on the IIA side at large volume. This is the desired shift-symmetric structure we want to start from.

2.2 Bulk Higgs vs. intersection-curve Higgs

Promoting our D7-brane to an $SU(6)$ stack, ζ becomes an $SU(6)$ adjoint. After further symmetry breaking to $SU(5)$ (or directly to the Standard Model), it can then contain the $\mathbf{5} + \bar{\mathbf{5}}$ (or $\mathbf{2} + \bar{\mathbf{2}}$) Higgses, with a shift symmetric Kähler potential. Models of this type have recently been considered in F-theory [37]. For reviews of model building in IIB/F-theory see e.g. [38].

Given that in type IIB and F-theory only the Kähler moduli T are *not* stabilized by fluxes in a supersymmetric way, we now assume that supersymmetry breaking is dominated by the F-terms of T. However, our bulk Higgs from D7-brane scalars has the peculiar feature that there is no Kähler moduli dependence of the Kähler metric [32,39]. Hence, assuming also shift symmetry and applying the supergravity formulae of [18], the normalized soft Higgs masses are simply given by

$$m_1^2 = m_2^2 = m_3^2 = m_{3/2}^2 \,. \tag{19}$$

The mass of the heavy Higgs doublet is therefore set to $m_S^2 = 2m_{3/2}^2$.

The Higgs mass matrix above is precisely of the form conjectured in a related way in [16]. However, here it follows with much less model dependence than expected earlier from a Giudice-Masiero-type analysis based on higher-dimensional operators and F_T. In a sense, this is our main 'success story' concerning a framework with a natural stringy shift symmetry. The remainder of this section is devoted to the related D7-brane intersection-curve Higgs where, as it turns out, the picture concerning a shift symmetry (and a possible prediction of the Higgs mass matrix in general) is much more complicated.

Most naively, the intersection-curve Higgs in type IIB orientifolds comes from the intersection locus of a $U(1)$ brane and a $U(2)$ brane stack. The two doublets H_u and H_d would in this case be two 4d zero modes of a 6d hypermultiplet living on the intersection curve of the D7-brane-stacks. This has a straightforward generalization to the $SU(5)$ GUT case, where a $\mathbf{5}$ and $\bar{\mathbf{5}}$ Higgs can originate in an analogous manner from a $U(1)$ and an $SU(5)$ stack (or, even more generally, from the intersection locus of a D_1 and the GUT divisor). For such matter fields, it is known that the Kähler potential has a Kähler-moduli-dependent prefactor [39]. We will now describe a localization process turning the previously discussed bulk Higgs into the intersection-curve Higgs. We aim in particular at understanding what happens to the shift symmetry.

To keep things as simple as possible, we ignore any possible GUT-superstructure. Instead, we start from a $U(3)$ D7-brane stack and deform it to the $U(1) \times U(2)$ stack. In this way, a 'brane Higgs' at the intersection curve can be continuously derived from a 'bulk Higgs' contained in the brane-deformation moduli ζ of the original $U(3)$ stack.

The relevant Kähler potential is that of (18), where we now assume that S and z are stabilized by fluxes (we will suppress the S- and z-dependence whenever possible) and

$\text{Im}(S) \gg 1$:

$$K = -\ln[-i(S - \overline{S}) - k_{D7}(z, \overline{z}, \zeta + \overline{\zeta})] \quad = \quad \frac{i}{(S - \overline{S})} k_{D7}(z, \overline{z}, \zeta + \overline{\zeta}) + \cdots$$

$$= \quad k_0 \, \text{tr}(\zeta + \overline{\zeta})^2 + \cdots . \tag{20}$$

Here we have first Taylor-expanded in k_{D7} (keeping only the linear term) and then restricted ourselves to a quadratic approximation in ζ absorbing the S- and z-dependence in the constant k_0.

Our ζ is in the adjoint of $U(3)$, containing in particular two doublets H_u and H_d in the complement of the $U(2) \times U(1)$ subgroup. So far, this is just group theory, but we go on to assume an actual deformation of our brane stack such that the surviving gauge group is $U(2) \times U(1)$ and the wave functions corresponding to the massless 4d fields H_u and H_d localize along the emerging intersection curve [40–42]. The survival of H_u and H_d as light 4d fields depends on geometrical data. More precisely, the internal wavefunctions of localised massless modes correspond to elements of the cohomology groups (see e.g. [41] for a review)

$$H^i(C, (L_a \otimes L_b^*)|_C \otimes K_C^{1/2}), \qquad i = 0, 1, \tag{21}$$

where C is the intersection curve, L_a and L_b denote the line bundles of the $U(2)$ and $U(1)$ branes whose curvature is the gauge flux F_a and F_b along the respective branes and K_C represents the canonical bundle of C. The modes corresponding to $i = 0$ and $i = 1$ characterize chiral $\mathcal{N} = 1$ multiplets in the representations $(\mathbf{2}_1, -1)$ and $(\mathbf{2}_{-1}, +1)$ under $U(2) \times U(1)$. Note that the difference of the number of such fields is given by the index $\int_C(F_a - F_b)$. In the sequel we assume that an appropriate choice of fluxes is made corresponding to $h^i(C, (L_a \otimes L_b^*)|_C \otimes K_C^{1/2}) = 1$ for $i = 0, 1$. One possible configuration would be e.g. a setup where C is a T^2 and the fluxes are chosen such that $(L_a \otimes L_b^*)|_C$ is the trivial bundle.

The relevant part of K is

$$K \quad \supset \quad k_0 |H_u + H_d^\dagger|^2 \,, \tag{22}$$

at least as long as the size of the $U(3)$-breaking deformation is sufficiently small. By slight abuse of notation, we denote the size of this deformation by $|\zeta|$. The 'smallness' of this quantity is defined as the requirement that the 8d wave functions of $H_{u,d}$ continue to be governed by the appropriate section $H^0(D, N_{D/X})$ of the normal bundle of the original $U(3)$ divisor D inside the Calabi-Yau X, such that the calculation of [32] continues to be valid. As already mentioned above, this will cease to be correct for sufficiently large values $|\zeta|$, when the relevant wave functions have become 6d wave functions described by (21) and supported only on the intersection curve.

In the case of torus orientifolds, the Kähler potential for D7-D7-brane matter is known [39] (based on [43, 44]). In particular, the kinetic term scales as $1/\sqrt{st}$, where s

and t are the real parts of S and T. Thus, focussing only on the H_u-kinetic term for brevity, we conjecture the general, qualitative form

$$K \supset \frac{k_0}{1 + |\zeta|^2 \sqrt{t/s}} |H_u|^2 \sim \frac{1}{s + |\zeta|^2 \sqrt{ts}} |H_u|^2 \,. \tag{23}$$

At $|\zeta| = 0$, this is obviously in agreement with what we inferred from [32] before. At $|\zeta|^2 \sim \sqrt{s/t}$, the typical distance between the $U(2)$ stack and the $U(1)$ stack reaches string length (one can convince oneself of this by returning to the original 10d string frame action and carefully thinking about the definition of the relevant 4d supergravity variables). Hence, at this value of $|\zeta|$, one expects a transition to the new scaling regime advertised earlier.

We give a more careful derivation of this transition in the Appendix. See also [45] for a discussion of the Kähler potential on matter curves.

Qualitatively (though not up to $\mathcal{O}(1)$ factors) the above also applies to the term $\sim |H_d|^2$. Unfortunately, we are unable to give a similar argument for the terms $\sim H_u H_d$ and $H_d^\dagger H_u^\dagger$. All we can say is that the coefficients of these terms will also change dramatically if $|\zeta|$ becomes so large that the Higgs fields localize on a curve. Furthermore, this change will necessarily involve s and t. Given this limited information, we parametrize the Higgs Kähler potential as

$$K \supset f_1(T, \overline{T})|H_d|^2 + f_2(T, \overline{T})|H_u|^2 + f_3(T, \overline{T}) H_u H_d + \text{h.c.} \,, \tag{24}$$

and we know that, in the shift-symmetric limit, these three functions tend to a single constant value, k_0 (cf. (23)). Everything beyond that appears to depend on the details of the geometry of the original $SU(3)$ surface and the emerging intersection curve. Note that f_1 and f_2, are, at least in principle, accessible by dimensional reduction at the two-derivative level. This is not obvious for f_3 since the latter does not contribute to the Higgs kinetic term before SUSY breaking. To approach the difficult issue of determining the functions f_i we next turn to a particularly simple setting.

2.3 Toy model with a torus intersection curve and implications for more general situations

Let us assume that our intersection curve is a T^2 and that the moduli space of the Calabi-Yau orientifold with its branes has a locus where the induced metric on this T^2 is flat. Furthermore, let us assume that within this locus there is a region where our T^2 degenerates to an S^1. This very specific simple case corresponds to the Higgs doublet pair being a hypermultiplet of a 5d gauge theory compactified on S^1. In this setting, and restricting attention only to the modulus governing the radius of this S^1, concrete statements can actually be found in the literature [46, 47]:

The 5d hypermultiplet possesses an $SU(2)_R$ symmetry which allows for a 'twisting' of the S^1-compactified theory by an element of this group. This is Scherck-Schwarz supersymmetry breaking and, moreover, it can be characterized in 4d $\mathcal{N} = 1$ language

as spontaneous SUSY breaking triggered by the F-term of the complexified S^1 volume modulus or radion superfield. By slight abuse of notation, we call this superfield T since it morally corresponds to the type IIB Kähler moduli discussed before. The 5d action can be displayed in 4d $\mathcal{N} = 1$ notation[5] as [47]

$$S = \int d^5 x \left\{ \int d^4\theta (T + \overline{T})(\underline{H}_u \underline{H}_u^\dagger + \underline{H}_d^\dagger \underline{H}_d) + \int d^2\theta \underline{H}_u \partial_5 \underline{H}_d + \text{h.c.} \right\} . \qquad (25)$$

Crucially, no term $\sim \underline{H}_u \underline{H}_d$ appears. Equivalently, as also worked out in [47], the scalar mass matrix induced by the $SU(2)_R$ twist is purely diagonal.

Thus, in the extremely simplified setting discussed above, a shift-symmetric structure does not arise. While an exchange symmetry appears natural (assuming that no gauge fluxes or other effects distinguish between the two complex scalars of our Higgs hypermultiplet), it can not be used to realize $\tan\beta = 1$. The reason is that off-diagonal terms in the mass matrix are completely missing. In other words, tuning for a zero eigenvalue corresponds to tuning for a vanishing mass matrix so that $\tan\beta$ remains undetermined.

This can be easily remedied by supplementing (25) with a supersymmetric mass term μ, which corresponds to the substitution $\partial_5 \to \partial_5 + \mu$ in (25). One then obtains a Higgs mass matrix [47]

$$m_1^2 = m_2^2 = |\mu|^2 + \frac{|F_T|^2}{(T + \overline{T})^2} \qquad \text{and} \qquad m_3^2 = \frac{2\mu F_T}{T + \overline{T}} , \qquad (26)$$

where the required tuning for a vanishing determinant, $\mu = \overline{F}_T / 2\text{Re}T$, leads to $\tan\beta = 1$ as desired. At the microscopic level, one can think of this μ term simply as of a 5d mass term or, more fundamentally, as of the vev of the A_5 or A_6 components in the underlying 5 or 6d gauged hypermultiplet theory.[6]

Alternatively, this mass term could remain zero at the classical level in the 5 or 6d theory, being generated by quantum effects in the 4d effective theory. For example, holomorphic μ-terms can be induced by suitable D3-brane instantons along holomorphic divisors [51], such that the size of the suppression of the coupling $\sim e^\tau H_u H_d$ depends on the specific regime in which the volume modulus τ of the instanton divisor is stabilized. This natural parametrical smallness may be advantageous.

Returning to the classical 5 or 6d analysis, we note that generating the μ term through A_5 or A_6 expectation values corresponds to twisting the theory by a $U(1)$. Any diagonal $U(1)$ symmetry under which the scalars of the chiral components of the hypermultiplet transform in a vector-like way (i.e. H_u and H_d have opposite charges) is in principle suitable. For example, such an extra $U(1) \subset SU(2)_R'$ appears in trivial dimensional reduction of 10d SYM to 6 or 5d as a subgroup of transverse $SO(4) \sim SU(2)_R \times SU(2)_R'$ rotations. However, for $SU(2)_R \times SU(2)_R'$ to be a symmetry of a hypermultiplet, the latter

---

[5] Of course, field normalizations have to be adjusted: $H_u = \underline{H}_u \sqrt{2\pi R(T + \overline{T})}$ and analogously for H_d.

[6] Recall that in 6d an $\mathcal{N} = 2$ hypermultiplet can not have a mass term (see [48,49] and the discussion in Sect. 5 of [50]).

has to be in a real representation of the gauge group (cf. [52] and Sect. 12.6 of [53]). In our context this presumably means that we would need to extend the Higgs field content.

Let us now, based on the intuition gained from the oversimplified torus or even 5d case, return to a more generic setting. In other words, the two Higgs doublets now come from a hypermultiplet in a 6d model, compactified on some Riemann surface. In this case, the twisting is by $SU(2)_R$, by the $U(1)_T$ structure group of the tangent bundle, and by the gauge groups (including $U(1)$s) of the intersecting brane stacks. As long as the D7-branes are holomorphic divisors of a Calabi-Yau orientifold, the $SU(2)_R \times U(1)_T$ twisting leaves an $\mathcal{N} = 1$ supersymmetry intact. By analogy to 5d radion mediation, one expects that SUSY-breaking by the F-terms of the Kähler moduli is equivalent to non-supersymmetric $SU(2)_R \times U(1)_T$ twisting. However, since the hypermultiplet scalars are singlets under the $U(1)_T$ structure group, this will leave us with an effective $SU(2)_R$ twisting in the scalar sector. This is similar to what we just saw in the 5d model of [47]. We are thus left without an off-diagonal mass term. As before, this can be cured by inducing a μ (and hence, as in (26), a $B\mu$ term) by either non-perturbative effects or, classically, through twisting by gauge-theory $U(1)$s. This twisting might be associated with D7-brane Wilson lines or with D7-brane-fluxes. In the latter case, parametrical smallness of the μ term might be difficult to achieve.

At a more fundamental level, the non-supersymmetric twists should be related to 3-form fluxes in the Calabi-Yau or, more precisely, to the warping induced by those fluxes [54]. Thus, the exact form of the function $f_i(T, \overline{T})$ should be indirectly accessible through dimensional reduction in the presence of warping [55]. In particular, it should be possible to understand in some detail how the Higgs mass matrix can be tuned by the flux choice. It would be very interesting but goes beyond the scope of this paper to actually perform such a calculation.

To sum up, our best choice for finding a shift symmetry in a Higgs sector on D7-branes appears to be associated with a pure *bulk* Higgs. In this case, Kähler moduli do not enter the Higgs-sector Kähler potential and a simple mass matrix determined entirely by $m_{3/2}$ results. By contrast, looking for an exchange symmetry appears to be more promising in situations with an *intersection-curve* Higgs. While the general analysis is involved, the case with a T^2 intersection curve and no significant gauge flux effect seems to allow for the desired structure. The intermediate regime, where the Higgs is in transition from an 8d to a 6d field, is even harder to analyze and we were only able to give very qualitative results.

3 $\mathcal{N} = 2$ D-term effects

Let us now pursue an alternative idea for realizing $\tan \beta = 1$, based on D-term effects within the context of an $\mathcal{N} = 2$ Super-Yang Mills theory. Our discussion is inspired by the 'D6-branes-at-angles' mechanism for SUSY breaking with $\tan \beta = 1$ of [22, 56]. We will attempt to formulate the central idea more generally and return to this particular scenario shortly.

To begin, we recall that in a UV completion of the Standard Model any extended supersymmetry has to be broken at least at the compactification scale m_C because of chirality.[7] However certain non-chiral subsectors, such as the Higgs or the gauge sector, may continue to enjoy an approximate $\mathcal{N} = 2$ symmetry below m_C. By contrast, it is also possible that breaking to $\mathcal{N} = 0$ occurs directly at the compactification scale. Thus, we see that a clear separation between the various breaking scales from the higher-dimensional extended supersymmetry all the way down to the non-supersymmetric SM can not be taken for granted. For example in toroidal compactifications of type IIA (see [57] and references therein), gauge theories on D6 branes are locally $\mathcal{N} = 4$, matter lives at points with $\mathcal{N} = 1$ or (real) curves with $\mathcal{N} = 2$, and the overall supersymmetry of the brane/o-plane network may be $\mathcal{N} = 1$ [58, 59] or even $\mathcal{N} = 0$ [60–62]. This can have interesting implications for the Higgs mass matrix and quartic potential.

3.1 Constraints on extended SUSY

Before coming to possible uses of extended SUSY, let us point out a crucial constraint within our setting: We are interested in situations where $(H_u, H_d^\dagger)$ form a hypermultiplet charged under an $\mathcal{N} = 2$ U(1) gauge theory. Recall that the $\mathcal{N} = 2$ D-term potential [48] (also dubbed 'P-term potential' [63]) arises from the terms

$$\mathcal{L} \supset \cdots + \frac{1}{2}\vec{P}^2 + g\phi^A \vec{P} \cdot \vec{\sigma}_A{}^B \phi_B^\dagger + \cdots \tag{27}$$

of the complete lagrangian. Here $\vec{P}$ is the SU(2)$_R$-triplet auxiliary field and $\{\phi^A\} \equiv (H_u, H_d^\dagger)$ are the scalars of the hypermultiplet. The scalar potential hence reads

$$V_{D, \mathcal{N}=2} = \frac{1}{2}\vec{P}^2 \,, \tag{28}$$

where

$$P_3 = g(|H_u|^2 - |H_d|^2) \tag{29}$$

and P_1, P_2 follow from P_3 by appropriate SU(2)$_R$ rotations of the SU(2)$_R$ doublet $(H_u, H_d^\dagger)$. From a 4D $\mathcal{N} = 1$ perspective, $P_1 + iP_2$ is an F-term while P_3 is a D-term. Thus,

$$V_{D, \mathcal{N}=2} = V_F + V_D = \frac{1}{2}|P_1 + iP_2|^2 + \frac{1}{2}P_3^3 \,. \tag{30}$$

Explicitly, the potential takes the form

$$
\begin{aligned}
V_{D, \mathcal{N}=2} &= \frac{1}{2}(-g\phi_\alpha^A \sigma^a{}_A{}^B \phi_B^{\dagger\alpha})(-g\phi_\beta^C \sigma^a{}_C{}^D \phi_D^{\dagger\beta}) \\
&= \frac{g^2}{2}\Big((H_u H_d + H_d^\dagger H_u^\dagger)^2 - (H_u H_d - H_d^\dagger H_u^\dagger)^2 + (|H_u|^2 - |H_d|^2)^2\Big) \\
&= \frac{g^2}{2}\Big((|H_u|^2 - |H_d|^2)^2 + 4|H_u H_d|^2\Big) \,.
\end{aligned} \tag{31}
$$

---

[7] A model with an $\mathcal{N} = 2$ breaking scale parametrically below m_C is impossible since it would have to contain massive chiral fermions to pair up with those of the SM.

Now we recall that our main interest in $\tan\beta = 1$ was the implication $\lambda = 0$, i.e., the flatness of the scalar potential for the light SM Higgs boson. This, however, is violated by $|P_1|^2 + |P_2|^2$. The P_3^2 term corresponds to the D-term potential

$$V = \frac{g^2}{2}(|H_u|^2 - |H_d|^2)^2 \,, \tag{32}$$

which arises in this form since the 4D chiral multiplets H_u and H_d are in mutually complex conjugate representations of the gauge group.

It is thus necessary to have all D- and F-terms which are not of the type (32) decouple from the effective theory in an approximately supersymmetric fashion. This is tantamount to breaking to an MSSM-like gauge sector with $\mathcal{N} = 1$ SUSY. In other words, it appears to be necessary to have at least some small hierarchy between the scales where the gauge sector breaks to $\mathcal{N} = 1$ and to $\mathcal{N} = 0$. Note that we talk here about the quartic potential of the massless eigenstate. The potential (31) alone as a function of H_u and H_d has charge-violating flat directions of the type $\langle H_u^0 \rangle = -\langle H_d^- \rangle = c$, which are however not preserved once the $SU(2)_L$-D-Term is included.

3.2   Decoupling the 'wrong' D-term potentials

The mechanism of decoupling (and thus the appearance of flat directions in the low energy theory) can be understood from a simple 4D toy model describing the extended SUSY sector in terms of superfields. It includes a chiral singlet coupling to the Higgs doublets

$$\mathcal{W} = \kappa S H_u H_d + \frac{1}{2}M S^2 + \dots \tag{33}$$

where κ can be related to the gauge coupling in the previous section via $\kappa = \sqrt{2}g$. The quartic Higgs coupling in such a simple model has for example been discussed in [11]. In our case, we are especially interested in the interpretation $\sqrt{2}S = A_4 + i\Phi_7$, appropriate to an $\mathcal{N} = 4$ D6-brane theory with internal gauge components A_4, A_5, A_6 and adjoint scalars Φ_7, Φ_8, Φ_9. We can then identify the chiral auxiliary field with two components of $\vec{P}$, e.g. $F^S \sim P_1 + iP_2$. Above the scale M, the singlet and Higgs F-term contributions to the scalar potential are

$$F^{S\dagger}F^S \;\; = \;\; \kappa^2|H_uH_d|^2 + M^2 S^\dagger S + \kappa M(H_uH_d S^\dagger + H_d^\dagger H_u^\dagger S) \tag{34}$$

and

$$F^u F^{u\dagger} + F^{d\dagger}F^d \;\; = \;\; \kappa^2 S^\dagger S(|H_u|^2 + |H_d|^2)\,. \tag{35}$$

This potential contributes to the quartic Higgs coupling and can modify the tree-level prediction for λ. At the scale M, the entire chiral multiplet S is integrated out in an $\mathcal{N} = 1$ supersymmetric fashion. The corresponding Feynman graphs are shown in Figure 1. This can for example be the compactification scale. Since we expect the effective theory below $\Lambda = M$ to be a consistent supersymmetric theory at the renormalizable level, the F-term potential must vanish once the multiplet is integrated out. At tree-level, this
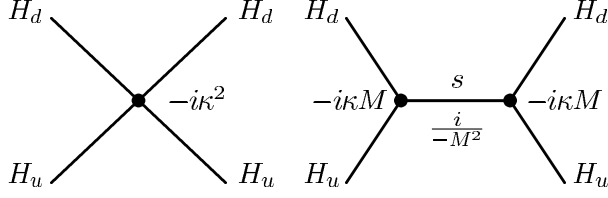
Figure 1: Due to destructive interference between these diagrams, the quartic F-term potential decouples at tree-level when the scalar s is integrated out. This decoupling is not exact if the scalar mass receives additional soft breaking contributions.

happens via the exchange of the massive scalar in S via the interaction in the third term of (34). Indeed, when expanding the graphs in Figure 1, one finds that the amplitude is of order $\mathcal{O}(p^2/M^2)$. There is therefore no dimension four quartic Higgs coupling in the effective theory[8] below $\Lambda = M$. This decoupling of the F_S potential is not exact any more once s receives a soft mass m_s^2, and the potential resulting from the matching at treelevel to the theory without the scalar multiplet reads

$$V_{\Lambda=M} = \kappa^2 \frac{m_s^2}{m_s^2 + M^2} |H_u H_d|^2 \overset{M \gg |m_s|}{\sim} \kappa^2 \frac{m_s^2}{M^2} |H_u H_d|^2 \, . \tag{36}$$

As one would expect, the quartic coupling still vanishes for $M/m_s \to \infty$. However, for moderate hierarchies between the overall soft scale and M, which could for example be the compactification scale (but also much lower), noticable corrections to $\lambda(m_S)$ and thus m_h can appear. In models where no clear hierarchy between breaking to $\mathcal{N} = 1$ and to $\mathcal{N} = 0$ is present, this has to be taken into account. For example, if the SM $SU(2)$ gauge theory originates form a brane theory with extended SUSY, a direct breaking to $\mathcal{N} = 0$ will lead to a significant quartic coupling $\sim g_2^2$ in the Higgs sector *without* any flat direction.

3.3 Using extended SUSY D-terms for the Higgs mass matrix

Naively, one would conclude from the above that all gauge theories under which our hyper is charged have to be broken to the 'surviving' $\mathcal{N} = 1$ SUSY much above the 'last' SUSY breaking scale. In this case, the quartic potential is flat along the direction of the SM Higgs if $\tan \beta = 1$ can be realized.

If, at the same time, we want to use a non-zero (field-dependent) FI term of such a gauge theory to generate a contribution to the Higgs mass matrix,

$$\mathcal{L} \supset g^2 \left(\xi + |H_u|^2 - |H_d|^2 \right)^2 \, , \tag{37}$$

---

[8] This is of course a well-known property of SUSY theories, and essential e.g. for SUSY GUT phenomenology. The same mechanism ensures that the divergent F- and D-term contributions of Kaluza-Klein modes decouple in the low-energy limit of extra dimensional models [64].

one encounters a problem: One finds opposite contributions to the soft masses of H_u and H_d, ruining $\tan\beta = 1$. Thus, $\xi = 0$ appears to be required and no useful effect of our extended-SUSY D-terms is left.

However, since we are here *not* dealing with a SM gauge group, we can try to avoid such a pessimistic conclusion by considering very small g. (Formally, this would be the limit $g \to 0$ with $g^2\xi$ fixed.) In this case, the induced quartic potential vanishes but the mass correction from ξ survives. The vanishing of the quartic potential allows us to consider an $\mathcal{N} = 2\ U(1)$ theory, which in turn gives us the choice between three FI terms. They are related by an $SU(2)_R$ rotation, which of course also affects our Higgs scalars. An interesting choice is to set the 'usual' FI term (the one associated to P_3) to a non-zero value, but working in a frame in which the hyper has the form

$$\{\Phi^1, \Phi^2\} \equiv \{(H_u - H_d^\dagger), (H_d^\dagger + H_u)\}\,. \tag{38}$$

Then, in analogy to (37), we find

$$\mathcal{L} \supset g^2 \left(\xi + |H_u - H_d^\dagger|^2 - |H_d^\dagger + H_u|^2\right)^2 \supset -4g^2\xi H_u H_d + \text{h.c.}\,, \tag{39}$$

which gives us a $B\mu$ term from a D-term. This is our interpretation of the very interesting suggestion for realizing the exchange symmetry (and hence $\tan\beta = 1$) in the context of intersecting D6 branes in [22]. Let us now scrutinize this particular implementation of the field-theoretic mechanism described above in more detail:

We recall that, in this context, the Higgs hypermultiplet lives at the intersection of a $U(2)$ stack of D6 branes (D6$_a$) with a single D6 brane (D6$_b$), where for definiteness we are working in the context of Type IIA orientifolds with D6-branes instead of their T-dual Type IIB version. While we assume the $SU(2)$ inside $U(2)$ to be the group of SM weak interactions, we remain agnostic concerning the SM hypercharge $U(1)_Y$. The latter can be some linear combination of the two $U(1)$s just introduced (or even just the $U(1)$ from D6$_a$) and further $U(1)$ groups.

To obtain a hypermultiplet Higgs, D6$_a$ and D6$_b$ have to intersect (non-generically) in a real line rather than (generically) in a point. The role of the $\mathcal{N} = 2\ U(1)$ theory can then be played by the gauge theory on D6$_b$, with ξ related in the standard way to a small deviation of the intersection angles from the SUSY condition. The smallness of g corresponds to a large volume of D6$_b$.

We now want to break the supersymmetry of the SM $SU(2)$ brane stack D6$_a$ in the appropriate way. Before doing so, let us first recall the situation in a standard $\mathcal{N} = 1$ compactification with $\mathcal{N} = 2$ Higgs sector. Without loss of generality, we start with a brane stack D6$_a$ filling the 468 plane, and choose a complex structure by defining that

$$(z_1, z_2, z_3) = (x^4 + ix^5, x^6 + ix^7, x^8 + ix^9)\,. \tag{40}$$

In flat space, we can think of z_i as transforming like a fundamental **3** of the "Calabi-Yau like" holonomy group $SU(3) \subset SO(6)$. As long as we remain in the standard $\mathcal{N} = 1$ picture, all D6/O6 positions are related to D6$_a$ by such $SU(3)$ rotations. For example,
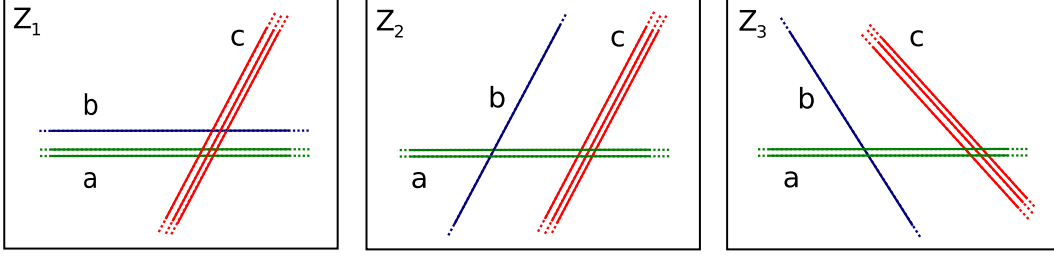
Figure 2: Sketch of the local intersection structure of the three relevant brane stacks in the D6-brane-realization of an exchange-symmetric Higgs mass matrix.

we can obtain D6$_b$ from D6$_a$ via

$$T(M_b) = \text{diag}(1, e^{i\theta}, e^{-i\theta}) \in SU(2) \subset SU(3). \tag{41}$$

Here, we denote by M_i the $SO(6) \subset SO(1,9)$ rotations in the spinor representation while $T(M_i)$ is the corresponding representation of $U(3) \subset SO(6)$ acting on (40). For our purposes, it is sufficient to supplement D6$_a$ and D6$_b$ with a (SM color) $SU(3)_c$ stack[9] D6$_c$, which has a chiral intersection with D6$_a$, as shown in Fig. 2. It is obtained from D6$_a$ via a rotation

$$T(M_c) = \text{diag}(e^{i\theta_1}, e^{i\theta_2}, e^{i\theta_3}) \tag{42}$$

with $\theta_1 + \theta_2 + \theta_3 = 0 \mod 2\pi$ as usual. As a result, the D6$_{a,b}$ system enjoys $\mathcal{N} = 2$, while the D6$_{a,c}$ system only enjoys $\mathcal{N} = 1$ supersymmetry.[10]

The position of D6$_b$ is characterized by two angles θ_2 and θ_3 in the z_2- and z_3-plane. As long as $\theta_2 + \theta_3 = 0$, this is just a parametrization of the $SU(2)$ rotation introduced above. We now go beyond $SU(2)$ to $U(2) \not\subset SU(3)$ by allowing for a small non-zero $\delta \equiv \theta_2 + \theta_3$, characterizing supersymmetry breaking from $\mathcal{N} = 2$ to $\mathcal{N} = 0$. It is well known [56,61] that, if we separate D6$_a$ and D6$_b$ along $\text{Im}(z_1)$ to avoid a tachyonic instability, then the mass matrix for the fields $\{\Phi^1, \Phi^2\}$ at the real intersection curve takes the form[11]

$$\begin{pmatrix} \mu^2 + \alpha\delta & 0 \\ 0 & \mu^2 - \alpha\delta \end{pmatrix}. \tag{43}$$

We would, by contrast, like to obtain the form

$$\begin{pmatrix} \mu^2 & \alpha\delta \\ \alpha\delta & \mu^2 \end{pmatrix}, \tag{44}$$

---

[9] Actually, the $SU(3)$ stack is introduced just as a simple illustration of the required SUSY breaking to $\mathcal{N} = 1$ of D6$_a$. What is crucial for all that follows is a reduction of the spectrum on D6$_a$ to that of an $\mathcal{N} = 1$ gauge theory. This is not realized by the $SU(3)$ stack, but by intersection with an O-plane. We assume that this latter intersection breaks SUSY in the same way as the $SU(3)$ stack and keep talking about the $SU(3)$ stack as a simple way of specifying a particular $\mathcal{N} = 1$ subalgebra.

[10] For a general discussion of the relative supersymmetry preserved by branes at angles we refer to [66].

[11] Other effects of D-term breaking on the spectrum in models with extended relative supersymmetries have been discussed recently in [65].

which corresponds to a rotation of (43) by

$$U = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & 1 \\ -1 & 1 \end{pmatrix} \in SU(2)_R \,. \tag{45}$$

This last rotation is automatically realized if we manage to ensure that $\{\Phi^1, \Phi^2\} = \{(H_u - H_d^\dagger), (H_d^\dagger + H_u)\}/\sqrt{2}$, with H_u, H_d being the scalars of the superfields corresponding to the $\mathcal{N} = 1$ SUSY respected by the D6$_{a,c}$ system. Thus, we need to 'rotate' the SUSY breaking induced by D6$_c$ on the D6$_a$ gauge theory by an $SU(2)_R$ twist U. All we would then have to do is to replace D6$_c$ by D6$'_c$, where the position of the latter follows from D6$_a$ via a rotation by U

$$M_c' \equiv U M_c U^\dagger \,, \tag{46}$$

where all objects are now matrices acting on the spinor representation of $SO(1,9)$. What is the appropriate choice for U? The condition for the unbroken supersymmetry parameter in the $D6_{a,c}$ system is

$$[\Gamma_{D_a}, M_c]\epsilon = 0, \tag{47}$$

where $\Gamma_{D_a} = \prod_i e_i^M \Gamma_M$ is defined in the usual fashion [66], and i runs over the brane-parallel directions. We want to implement U such that the new condition for the unbroken parameter ϵ' of the $D6_{a,c'}$ system,

$$[\Gamma_{D_a}, M_c']\epsilon' = 0, \tag{48}$$

automatically holds for $\epsilon' = U\epsilon$. Furthermore, ϵ' should precisely correspond to the $\mathcal{N} = 1$ supersymmetry in which the hypermultiplet takes the desired form (38). The position of the D6$_a$ stack in the 468 direction as depicted in Figure 2 corresponds to

$$\Gamma_{D_a} = \Gamma_4 \Gamma_6 \Gamma_8 \,. \tag{49}$$

The real line in the $D_{a,b}$ system on which the hypermultiplet resides was obtained by setting $\theta_1 = 0$. This determines the two supersymmetries of the Higgs system. To construct U in the spinor representation, we now consider the $SO(4) \subset SO(6)$ which rotates the coordinates 6789 while leaving the real line of the Higgs system untouched. The generators of this $SO(4) \cong SU(2) \times SU(2)_R$ can be represented as

$$T_1 = \frac{i}{8}\left([\Gamma^6, \Gamma^9] - [\Gamma^7, \Gamma^8]\right), T_2 = \frac{i}{8}\left([\Gamma^6, \Gamma^8] - [\Gamma^9, \Gamma^7]\right), T_3 = \frac{i}{8}\left([\Gamma^6, \Gamma^7] - [\Gamma^8, \Gamma^9]\right)$$
$$T_1^R = \frac{i}{8}\left([\Gamma^9, \Gamma^6] + [\Gamma^8, \Gamma^7]\right), T_2^R = \frac{i}{8}\left([\Gamma^6, \Gamma^8] + [\Gamma^9, \Gamma^7]\right), T_3^R = \frac{i}{8}\left([\Gamma^6, \Gamma^7] + [\Gamma^8, \Gamma^9]\right).$$
$$\tag{50}$$

The subgroup $SU(2)$ is contained in the $SU(3)$ "holonomy" group and leaves the unbroken supersymmetry invariant. The subgroup $SU(2)_R$ however does not respect the complex structure, and we can use it to rotate the unbroken supersymmetry (note however that arbitrary $SU(2)_R$ transformations on M_c can break supersymmetry in the D6$_{a,c'}$ system completely). The generator T_3^R performs a "factorizable" rotation and therefore

19

commutes with M_c. It leaves M_c, and consequently the unbroken supersymmetry, invariant. The generator T_2^R on the other hand satisfies $[T_2^R, \Gamma_{Da}] = 0$ since it leaves the branes in D6$_a$ manifestly invariant. As a consequence, for the choice

$$U = \exp[i\frac{\pi}{2} T_2^R] \tag{51}$$

the SUSY condition for the D6$_{a,c'}$ system becomes

$$0 \stackrel{!}{=} [\Gamma_{Da}, M_c']\epsilon' = [\Gamma_{Da}, U M_c U^\dagger]\epsilon' = U[\Gamma_{Da}, M_c]U^\dagger \epsilon' \tag{52}$$

which is satisfied for $\epsilon' = U\epsilon$. U now rotates the two unbroken supersymmetry parameters of the Higgs sector into each other (by an angle of $\pi/2$). As a result, the $\mathcal{N} = 1$ supersymmetry of D6$_a$ respected by the intersection with D6$_c'$ is precisely the one in which the $\mathcal{N} = 1$ superfields H_u and H_d receive a $B\mu$ term mass correction $\sim \delta$ according to (39).

A different perspective might also be useful: Let us stick with M as the rotation of D6$_a$ to D6$_c$ and denote by $(1 + \delta)N$ the rotation of D6$_a$ to D6$_b$. Here δ is now a small matrix, such that $N \in SU(2)$ and $(1 + \delta) \in U(2)$. Rather than rotating M by U, we instead rotate $(1 + \delta)N$ by $U^\dagger$, which is clearly equivalent. We thus have

$$(1 + \delta)N \to U^\dagger(1 + \delta)NU = (1 + U^\dagger \delta U)N. \tag{53}$$

The last equality follows since $N \in SU(2)$ and $U \in SU(2)_R$, with $SU(2)$ and $SU(2)_R$ the two commuting factors of $SO(4)$ or, equivalently, the upper-left and lower-right 2×2 blocks of $SU(4)$. We now see very clearly that one can modify the SUSY-breaking rotation $(1 + \delta)$ such that it induces an F-term instead of a D-term. This is not what is usually done in SUSY-breaking by angles, since one usually just modifies angles in the factorized geometry of three complex planes. But this is clearly not the generic possibility of rotating branes, and the appearance of an F-term corresponds to a general rotation after which the D6-brane ceases to be a lagrangian brane. In a Type IIB dual such an F-term translates into F-term breaking by non-holomorphicity of the branes and/or by suitable gauge fluxes. Note that the latter option requires gauge fluxes which are not the pullback of ambient space forms and is thus not available in simple toroidal geometries.

In summary, we have seen that deriving a $B\mu$ term from the D-terms of extended supersymmetry is field theoretically possible, but a detailed understanding of the hierarchies of SUSY breaking is mandatory. In particular, in order to realize an exchange symmetric Higgs mass matrix, one has to make sure that the soft diagonal contributions are strongly suppressed. For example, any model that exploits $\tan\beta = 1$ to predict approximately vanishing quartic coupling but generates $B\mu$ via F-term contributions must ensure that the same F-term decouples sufficiently from the Higgs potential below m_S. This might be possible in the limit where the effective gauge coupling of the gauge theory from which this F-term originates is small. Finding an explicit stringy realization with D6 branes on tori appears doable according to the discussion given above. We stress, however, that the torus geometry, which is crucial for the viability of the $SU(2)_R$ symmetry underlying the argument, is at odds with a strong hierarchy between the gauge coupling associated with the theory from the which the F-term results and the Standard

Model: Due to the simple cycle structure on T^6 the only way to ensure strongly differing gauge couplings and thus brane volumes is to include branes with high wrapping numbers, which, on the other hand, is constrained by tadpole cancellation conditions. More interesting would be an understanding of this mechanism in the context of a proper Calabi-Yau geometry.

4 Corrections to the Quartic Coupling

4.1 Radiative Corrections to the Quartic Coupling

The dominant radiative corrections to the quartic coupling in the UV come from two sources: shift/exchange symmetry violating corrections to the Higgs mass matrix will lead to a deviation from $\tan\beta = 1$ and thus indirectly to corrections of the effective quartic coupling in the SM. Furthermore, the quartic coupling itself receives threshold and decoupling contributions when the heavy Higgs states, sfermions (in particular the top partners) and higgsinos and gauginos are integrated out near m_S. In unbroken SUSY, the relation (1) between the D-term quartic coupling and the gauge couplings generally only holds to higher orders in manifestly supersymmetric renormalization schemes like $\overline{DR}$. We find that the corresponding conversion of the quartic coupling to $\overline{MS}$ yields a numerically negligible contribution for our purposes. Our discussion of radiative corrections to the Higgs mass in high-scale supersymmetry is somewhat analogous to the one performed in [15]. However, while these authors have concentrated on the case $\cos^2 2\beta \sim 1$ which is now disfavored by data, we treat the complementary case $\tan\beta \sim 1$.

Corrections to the Higgs Mass Matrix: In a previous paper [16], we have studied the first effect in order to quantify the impact of shift symmetry violating interactions on the observed Higgs mass. The starting point is the usual mass matrix with universal entries (5) from which one obtains $\tan\beta = 1$ and thus $\lambda(m_S) = 0$ from the tree level D-term potential (see Section 1.1). Radiative corrections, in particular from the superpotential Yukawa couplings to the top quark, necessarily modify this texture at 1-loop, leading to a deviation from $\cos^2 2\beta = 0$. For shift symmetric scenarios in which the symmetry can be assumed to be exact in the uncompactified higher dimensional theory, there are two related corrections to the mass matrix. Since the transition from the MSSM to the SM is not at the compactification scale m_C, but rather at a lower scale m_S, we can collect the effects from the shift symmetry violating interactions by the RGE running from m_C down to m_S[12]. The corrections are of the form [16]

$$\frac{\delta m_{Hi}^2}{m_{Hi}^2} \sim \frac{6\overline{y_t^2}}{16\pi^2} \log\left(\frac{m_S}{m_C}\right) \tag{54}$$

---

[12]The same strategy can be applied to estimate the corrections to the exchange symmetric scenarios. In [22], the authors estimate the corrections to the Higgs mass with a different approach, but essentially arrive at the same conclusions

and consequently, the shift violating (SV) correction to the quartic coupling is given by

$$\delta\lambda_{SV}(m_S) \sim C\, \frac{g_2^2 + g_1^2}{8} \left| \frac{6\overline{y_t^2}}{16\pi^2} \log\left(\frac{m_S}{m_C}\right) \right|^2. \tag{55}$$

where the constant C depends on how the electroweak scale is tuned light and was assumed to be $C = 2$ for definiteness. When integrating out $\tilde{t}_i$, we get decoupling contributions to the squared Higgs masses from quadratically divergent self energy graphs. At 1-loop, they include an ambiguity concerning the choice of matching scale and amount to an additive constant of $\mathcal{O}(1)$ to the log. For $m_C \gg m_S$, the constant is therefore negligible compared to the log and can be absorbed in a slight shift of m_S.

Integrating out heavy MSSM states: The second type of radiative corrections originating from the stops (T) is the one to the quartic coupling itself [67],

$$\delta\lambda_T(m_S) = \frac{3y_t^4}{16\pi^2} \left[\frac{X_t^2}{m_{\tilde{t}}^2}\left(1 - \frac{X_t^2}{12m_{\tilde{t}}^2}\right) + 2\log\left(\frac{m_{\tilde{t}}}{m_S}\right) \right] \tag{56}$$

where $X_t = A_t - \mu\cot\beta \approx A_t - \mu$ is the effective trilinear coupling. We have assumed negligible splitting among the top partners. The correction $\delta\lambda_T$ is nominally suppressed by one loop factor less than λ_{SV} in Eq. (55). The first term[13,14] in the brackets in Eq. (56) is the 1-loop effective A-term contribution from the finite graphs with four external SM Higgs fields. The second term encodes the threshold which arises if the stops are not integrated out at $m_{\tilde{t}}$. While the SM running of the quartic coupling at high scales becomes increasingly flat due to a cancellation between Yukawa and gauge contributions, and the overall sensitivity on m_S is thus suppressed, this is not necessarily true for the individual thresholds. Thus, in addition to the effects discussed in [23], we found that for moderate splittings $m_{\tilde{t}}/m_S$ and intermediate values of X_t, the effective A-term correction and thresholds in (56) are of similar size and can cancel.

Assuming that X_t and $m_{\tilde{t}}$ are roughly of the same order (splitting between $\tilde{t}_{L,R}$ is neglected here), the first contribution is maximized for $X_t^2 = 6m_{\tilde{t}}^2$. For $X_t^2 = 0\ldots 6m_{\tilde{t}}^2$ it is in the range

$$\delta\lambda_T(m_S) = 0\ldots 3 \times \frac{3y_t^4}{16\pi^2}. \tag{57}$$

The calculation given in [16], where the focus was on effects from shift symmetry violation and the viability of the shift symmetric scenario, does not include these corrections. However, they should be taken into account together with the SUSY thresholds and the shift violating contributions to get a meaningful estimate of the value and uncertainty in m_S. The impact of these radiative corrections to λ on the physical Higgs mass is illustrated in Figure 3 for $\delta\lambda_{SV}$ and $\delta\lambda_T$. The shaded Higgs mass range corresponds to

---

[13]We thank Manuel Drees for reminding us of the potential importance and different quality of this contribution.

[14]Note that a negative correction to λ can be obtained from the 1-loop stop decoupling corrections for $X_t > \sqrt{12}m_{\tilde{t}}$. A nonperturbative regime can arise from extreme values of X_t, and one has to be careful not to introduce charge- or color breaking minima [68].
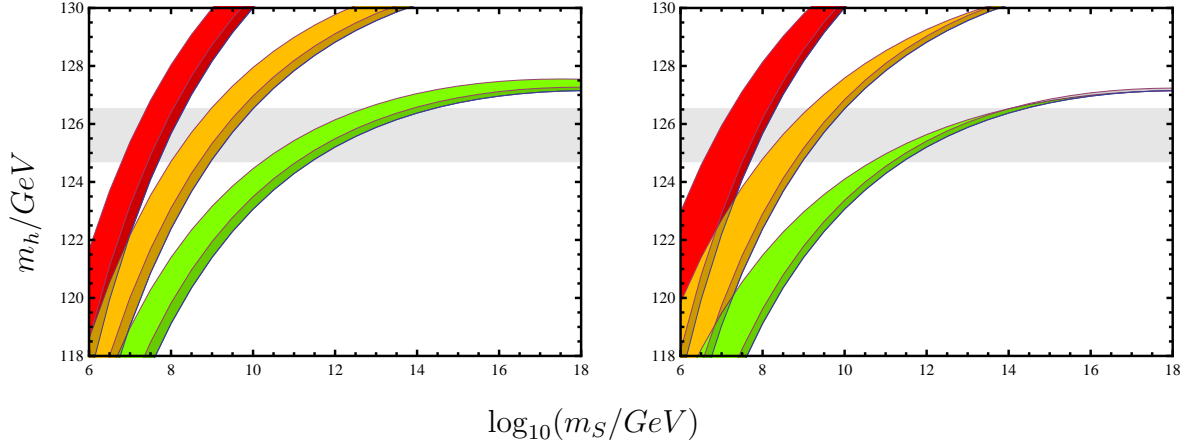
Figure 3: The impact of squark decoupling corrections to the quartic Higgs coupling (left) and shift/exchange symmetry violation (right) on the physical Higgs mass. The narrow dark(broad light) bands are for $X_t^2 = m_S^2 \, (6m_S^2)$ for the decoupling contributions from top partners, and $m_C = 10^2 \, m_S(\sqrt{m_S m_{Pl}})$ for the shift symmetry violation. The top quark masses are $m_t = 175.5, 173.5, 171.5$ from upper (red) to lower (green) band. The scale m_S should be understood as the effective SUSY scale as defined in (64).

the union of the ATLAS and CMS (1σ) errors quoted with the observation [1, 2], with statistics and systematics added in quadrature. The three colored bands correspond to the current central value given by PDG for the top quark mass ($173.5 \pm .6 \pm .8$ GeV, [69]) and $\pm 2\sigma$, again with errors added in quadrature. It is noteworthy that both the shift symmetry violating corrections and the threshold corrections to m_h owe their relative smallness to the RGE evolution of y_t towards smaller values in the UV with $y_t^4(10^9, 10^{16}, 10^{19} \text{ GeV}) \approx 1/9, 1/27, 1/40$.

Let us now consider the remaining threshold corrections in the limit $\tan\beta = 1$. They can be found e.g. in [11, 70]. The higgsino and gaugino threshold contributions are somewhat more complicated than the stop thresholds. We take the expressions given in [11] and apply some simplifications. First of all, we assume

$$r_1 \equiv \frac{M_1}{\mu} = r_2 \equiv \frac{M_2}{\mu} = \frac{M_\lambda}{\mu} \equiv r, \quad \tan\beta = 1, \quad \lambda^{LO} = 0, \quad (58)$$

where the higgsino mass μ should not be confused with the renormalization scale, and M_1, M_2 are the electroweak gaugino masses. The corrections to $\lambda(m_S)$ given in [11] then reduce to

$$\delta\lambda_{GH}(m_S) = \frac{\tilde{b}_\lambda}{16\pi^2} \left[\log \frac{\mu}{m_S} + \frac{(r-1)(r+1)^2 + 2(r-3)r^2 \log r}{2(r-1)^3} \right] \quad (59)$$

where

$$\tilde{b}_\lambda = \frac{1}{2}(-g_1^4 - 2g_1^2 g_2^2 - 3g_2^4). \quad (60)$$

23

In leading log approximation this becomes

$$\delta\lambda_{GH} \approx \frac{\tilde{b}_\lambda}{16\pi^2} \log \frac{m_\chi}{m_S} \tag{61}$$

where $m_\chi \equiv \max(\mu, M_\lambda)$. Note that we define g_1 in SM normalization rather than GUT normalization, and that our quartic coupling is normalized as $2\lambda(\text{this paper}) = \lambda([11])$.

Finally, there is a threshold from the heavy Higgs states, which will complete the MSSM-like 4D spectrum. We use the result in [70], again in the limit $\tan\beta = 1$, and find

$$\delta\lambda_A = -\frac{1}{16\pi^2}\frac{1}{4}\tilde{b}_\lambda \log \frac{m_A}{m_S} \tag{62}$$

where m_A is the mass of the heavy Higgs doublet, in our case $m_A^2 \sim 2B\mu$. The constant decoupling contributions from mixed heavy-light Higgs diagrams are found to vanish in the limit $\tan\beta = 1$. As a quick consistency check[15], note that the logs in $\delta\lambda_T + \delta\lambda_{GH} + \delta\lambda_A$ reproduce the leading log of the SM running of the quartic coupling. Specifically, in the $\lambda^{LO} = 0, \tan\beta = 1$ limit, we obtain

$$16\pi^2\frac{\partial}{\partial \log m_S}\delta\lambda_{T+GH+A} = \left[-\frac{3}{4}\tilde{b}_\lambda - 6y_t^4\right] = b_\lambda^{SM,1-loop}\Big|_{\lambda=0}. \tag{63}$$

which means that the quartic coupling obtained from the SUSY theory as a function of m_S "runs" like the SM quartic coupling to leading log precision. This lets the unphysical matching scale m_S drop out of our Higgs mass prediction (for small variations of m_S).

Ignoring the constant thresholds for now, one can find the value of the unphysical scale m_S for which the threshold logs cancel among each other. This "effective SUSY scale" to leading log precision, $m_S = m_S^{eff}$, is given by

$$m_S^{eff} = \left[m_A^{-\tilde{b}_\lambda/3} m_{\tilde{t}}^{8y_t^4} m_\chi^{4\tilde{b}_\lambda/3}\right]^{1/(\tilde{b}_\lambda + 8y_t^4)}. \tag{64}$$

Thus, in leading log approximation, this scale choice allows us to set the quartic coupling to the tree level relation[16] $\lambda(m_S^{eff}) = (g_1^2(m_S^{eff}) + g_2^2(m_S^{eff}))\cos^2 2\beta/8$. The denominator of the exponent in Eq. (64) corresponds to the SM β function of λ. For vanishing SM running, m_S is therefore undetermined in this approximation. Note that in our derivation of the scale m_S^{eff}, we have assumed that the constant thresholds can be neglected.

What is the relative importance of the threshold contributions compared to the effective A-term contribution? It turns out that the former can be negative and of similar magnitude as the latter even for the extreme case $X_t^2 = 6m_{\tilde{t}}^2$. For example, for $m_S = 10^9$ GeV, we have $y_t^4 \sim 0.11$ and $\tilde{b}_\lambda \sim -0.24$. We choose the matching scale at the mass scale of the heavy Higgs doublet, $m_S^2 = m_A^2 \sim 2B\mu$, thus resumming the log

---

[15]We find a discrepancy between the heavy Higgs thresholds given in [70] which we use here, and [11].

[16]Since we have used the approximation $\cos^2 2\beta = 0, \lambda^{LO} = 0$ in the threshold formulae, this is only consistent if $\cos^2 2\beta$ is of loop-suppressed size. This is generally the case in the shift/exchange symmetric models which we discuss in this paper.

in $\delta\lambda_A$. For definiteness, consider a scenario where the Higgs mass matrix has no soft mass contributions. This situation occurs in some models discussed in this paper. Then, $\mu^2 \sim B\mu$, and thus $\mu \sim m_S/\sqrt{2}$. For vanishing splittings $r = 1$ and $m_{\tilde{t}} \sim \mu$, we get $\delta\lambda_{GH} \sim \tilde{b}_\lambda/(4\pi)^2 \sim -0.0015$ and $\delta\lambda_{T,log} \sim -0.0015$. Although we have $\tilde{b}_\lambda \log(\mu/m_S) > 0$, the total GH contribution is negative because the constant part dominates over the log for $r \to 1$. Compare this to the "worst case" effective A-term contribution $\delta\lambda_{T,X_t} \sim 0.006$. To conclude, we find that even for relatively large $X_t^2 < 6m_{\tilde{t}}^2$, the radiative corrections to the effective quartic coupling $\lambda(m_S)$ can be negative, and can cancel the shift violating contributions in some scenarios. At higher soft breaking scales, the relative importance of the y_t^4 corrections will dwindle, and the gaugino-higgsino thresholds will dominate the corrections to $\lambda(m_S)$.

4.2 Models with $\lambda < 0$

As we have discussed, the SM running of the quartic coupling leads us into a regime where $\lambda < 0$ far below the Planck scale if the Higgs/top quark masses turn out to be respectively at the lower/upper end of their current experimentally allowed ranges. It is hence interesting to analyze whether a supersymmetric/stringy UV completion can be realized in this classically unstable regime.

The naive answer is 'no' since the MSSM D-term potential always implies a non-negative tree-level quartic coupling at the soft scale (cf. the familiar tree-level relation (1)). In fact, we have so far mainly exploited the limiting case $\tan\beta = 1$ and $\lambda(m_S) = 0$, thereby pushing the SUSY breaking scale as far up as possible. Of course, λ can receive loop corrections of either sign. Going beyond this, we will now appeal to classical corrections and argue that a *sizable* negative quartic coupling is a possibility that should be taken seriously.

It is well-known that the upper bounds on the mass of the light Higgs in the MSSM are loosened in extensions like the NMSSM, where the F-term potential of the singlet provides a contribution to the quartic coupling. Such effects can, of course, also be relevant in high-scale SUSY breaking. We will base our discussion, once again, on the toy model (33), also considered in [11]. In other words, we introduce a singlet chiral superfield S, allowing for a renormalizable coupling to the doublets and a supersymmetric mass term.

While the addition of a singlet scalar to the Higgs sector may seem ad hoc in field theory, it is quite natural in string models with soft breaking scales near the compactification scale. For example, if the chiral Higgs doublets H_u, H_d originate from a hypermultiplet on intersecting branes as proposed e.g. in [61], the bulk gauge theory degrees of freedom couple precisely in this manner[17] [71,72]. The degree of decoupling of the bulk F-terms from the 4D effective theory is then a crucial issue (see also section 3.1). We have seen that a non-vanishing quartic coupling will generally appear in the effective theory after integrating out the singlet scalar if the latter has a nonvanishing soft mass m_S.

---

[17]Such couplings to SM singlets can also occur if H_u and H_d originate from separate matter curves, in which case S is charged and can become massive via instanton effects.

We now consider situations where $\tan\beta = 1$ and, at the same time, $-M^2 < m_s^2 < 0$. As we will see, the main conclusion can be summarized in the equations

$$\mathcal{W} = \kappa S H_u H_d + \frac{M}{2}S^2 + \dots \quad \overset{\tan\beta=1 \; ; \; m_s^2<0}{\Longrightarrow} \quad V_{\Lambda<M} = \frac{\kappa^2 m_s^2}{M^2 + m_s^2}|H_0|^4 < 0 \,, \qquad (65)$$

where H_0 is the Standard Model Higgs doublet (which is massless at the high scale). The crucial point is that, in spite of the obvious tree-level instability, the running of λ will save the theory in the expansion around $H_0 = 0$ (until, of course, at very low energies, the negative mass squared of H_0 becomes relevant).

Let us now make this point in somewhat more detail, following the effective field theory from high to low energy scales. We will assume for simplicity that the Higgs mass entries $m_1^2 = m_2^2 = m_3^2$ are somewhat (but not hierarchically) smaller than M^2 and $-m_s^2$, which are of the same order of magnitude. The full classical scalar potential reads

$$\begin{aligned} V_{\Lambda>M} &= |\kappa H_u H_d + MS|^2 + |\kappa S H_u|^2 + |\kappa S H_d|^2 + \left\{ \kappa\overline{\mu}S(|H_u|^2 + |H_d|^2) + \text{h.c.} \right\} \\ &+ m_1^2|H_u + H_d^\dagger|^2 + \frac{g_2^2 + g_1^2}{8}(|H_u|^2 - |H_d|^2)^2 + \frac{g^2}{2}|H_u \epsilon H_d^\dagger|^2 + m_s^2|S|^2 \end{aligned} \qquad (66)$$

At energy scales high enough to neglect all masses, we have just the positive-definite quartic D and F term potentials and no flat directions are left. At smaller energies (smaller field vevs) one discovers an approximately flat direction. It corresponds to simultaneously switching on $H_{u,d}$ and S. In fact, including even higher powers in the analysis, one sees that that the potential falls slightly below zero when moving from zero into this flat direction. After integrating out the superfield S, this behavior translates to a negative treelevel quartic coupling for the light Higgs doublet H_0 in the effective theory below M (cf. (65)).

The potential at $H_0 = 0$ is unstable but we know (given the stability of our original softly broken SUSY model) that a stable minimum with all vevs non-zero exists. However, we propose to follow the RG flow of the theory near $H_0 = 0$. Our crucial point is that this is consistent and, once we reach the regime where $\lambda > 0$, we just continue to run conventionally further down until standard electroweak symmetry breaking sets in. Of course, the previously discovered negative-energy minimum is still present and 'our' minimum is only metastable. The situation is depicted in Fig 4.

Before worrying about the consistency of our analysis, let us first see what we would gain. Clearly, we can significantly loosen the upper bound for the heavy MSSM masses in our model, at least if the top quark turns out to be relatively heavy. For example, for $m_t = 173.5$ GeV, the 2-loop running gives us a minimum value of the SM quartic coupling in $\overline{MS}$ of $\lambda(10^{17.5}\text{ GeV}) \approx -0.016$. Neglecting the running of the potential between M and m_S, a negative correction of this size can be obtained for $\kappa \approx 1$ and $M \approx 8|m_s|$. Obviously, we can even shift the SUSY breaking scale all the way up to the Planck scale if we wish.

But is this picture consistent? While the potential is unstable around $H_0 = 0$ due to the negative effective quartic coupling, this instability is not tachyonic. There is no (quadratic) mass instability or even tadpole. Thus, we encounter no technical problem in
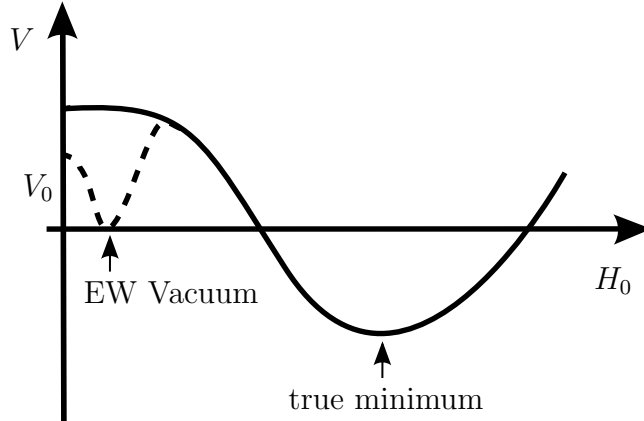
26

Figure 4: A sketch of the effective potential with negative quartic coupling before running (solid) and after running to low energies (dashed). An uplift term to ensure vanishing cosmological constant in the electroweak (EW) vacuum has been added.

doing perturbation theory (with a massless Higgs) around this extremum. The quartic instability introduces no mass scale into our model and hence, given a sufficiently smooth vacuum state, we can 'live' for a long time at $H_0 = 0$. The dynamics is governed by an effective field theory (with UV and IR cutoff), which can in principle be tested rather precisely (cf. Appendix B for more details).

This is, of course, in no way surprising from the perspective of cosmological inflation, where one is used to analysing potentials in perturbation theory around points with $V'' < 0$ (see, e.g., [73] for some early and more recent examples). If one also has $V' = 0$, the lifetime of some Hubble-sized patch on the 'top of the hill' is limited by the quantum fluctuations of the inflaton, which are controlled by the Hubble scale H. Obviously, there are also the familiar dS space IR divergences [74] and, in general, late-time non-perturbative effects [75]. However, all of this does not affect our application since, given the usual tuning of the cosmological constant, there is no high Hubble scale. Instead, as explained in the Appendix B, we can alway use a low IR cutoff and this cutoff sets the scale for both quantum diffusion and classical instability.

All we really need for our purposes is to calculate how λ, defined originally at a scale $\mu \sim M$, changes with μ. The definition of this λ is, very intuitively, via a 4-point correlation function at distances $\sim 1/\mu$ (hence without IR sensitivity) and with a UV cutoff Λ_{UV} not too far above μ. The answer to the question of the μ-dependence of λ is, of course, the conventional SM β-function. Thus, if we continue to calculate λ at lower and lower energy scales, its sign will eventually change and we return to the firm ground of perturbation theory around a stable extremum (again, obviously, before μ becomes comparable to the Higgs mass parameter).

While we believe that the arguments given above (and in App. B) are sufficiently convincing at the leading-log level, a more careful field theoretic discussion is certainly worthwhile. In particular, it is interesting to investigate how far one can push higher-orders perturbation theory in this classically unstable regime. Another interesting ques-

tion is how our universe has ended up in the radiatively generated, metastable state we just argued for. This may find a resolution along the lines suggested in [76] in a related context.

We reiterate that, as we have just seen, a stringy UV completion might occur far above the scale where λ turns negative. In the region with negative λ, we have no proper equilibrium field theory (in particular the Hamiltonian is unbounded below), but sufficient control in an effective theory with IR cutoff to answer 'short-time-scale' questions. Most importantly, we control the running in this regime.

5   Conclusions

The discovery of a Higgs-like boson with a mass of ~ 126 GeV has provided a measurement of the last undetermined parameter of the Standard Model: the quartic Higgs coupling λ. While it remains to be seen whether the minimal Higgs mechanism truly is the correct theory of electroweak symmetry breaking at and beyond the electroweak scale, all known properties of the new Higgs-like particle such as decay branching fractions and production rates are in agreement with SM predictions (see e.g. [77]). Absent new physics, the Higgs quartic coupling in the SM generically runs to $\lambda = 0$ and even negative values at some high renormalization scale $\mu_\lambda \gtrsim 10^8$ GeV for most of the experimentally allowed range of top quark and Higgs mass values (see e.g. [7–12]). For the purpose of this paper, we have taken this (and the current lack of evidence for new physics at colliders) as a hint that the scale μ_λ of vanishing quartic coupling might tell us something about the nature of the UV completion of the SM. This general idea is of course not new. There have for example been proposals of unified theories [19] and asymptotic safety scenarios [78] which have made predictions of the Higgs mass based on the UV boundary condition $\lambda(Q) = 0$ and $\lambda(Q) = \beta_\lambda(Q) = 0$ before the fact.

We assume in this paper that superstring/M-theory is the correct description of quantum gravity and fundamental interactions, and have identified models with high-scale supersymmetry (for a review of alternatives see e.g. [79], for a recent discussion of the phenomenology of the low intermediate regime see [80]) which naturally exhibit approximately flat directions in the Higgs potential, thus providing the UV boundary condition $\lambda(Q) = 0$ for the RG running (as well as $m_h^2(Q) = 0$ at tree level in the case of shift symmetric models). Most of our arguments hinge on the MSSM treelevel relation $\lambda(\tan\beta = 1) = 0$ which results from the flat directions of the D-term potential. In the decoupling limit in which we are working, the value $\tan\beta = 1$ is tied to a particular structure of the Higgs mass matrix where all entries are equal with $B\mu = m_1^2 = m_2^2$. Such mass matrices occur automatically if the Higgs sector features a shift symmetry [16] or an exchange symmetry [22]. In the latter case, only $m_1^2 = m_2^2$ is imposed by the symmetry, and $B\mu \approx m_i^2$ is a by-product of the tuning of the electroweak scale.

In our previous paper [16], we have identified shift symmetric Higgs sectors as a plausible explanation for $\tan\beta = 1$ at a high scale. While such shift symmetries have been known to appear in heterotic orbifold models for Wilson line type bulk fields, it is

28

unclear to what degree they are preserved when going from orbifolds to smooth Calabi-Yaus. In this paper, we have focused on type II/F-Theory compactifications with D6 and D7 branes. In analogy to the heterotic case, the obvious place where one would expect shift symmetries of the required type are again Wilson lines on D6 or D7 branes. It is important to note that the shift symmetry needs to be realized in the correct Kähler variables of the 4D effective supergravity theory in order to yield the desired soft masses. Not all components of the Higgs field can exhibit a shift symmetry simultaneously (this would forbid all terms in the Kähler potential). This, along with the appearance of a shift-violating Chern-Simons term, prevents the desired shift symmetry from being generically obeyed by D7 Wilson lines. In the case on D6 branes, the complex variables of the supersymmetric theory each combine one real Wilson line degree of freedom with one brane scalar describing normal movement. We have argued that such bulk Higgs constructions and their type IIB duals (with a Higgs from D7 brane scalars) indeed exhibit shift symmetries of the required type.

A further possibility to implement shift/exchange symmetric models might be given by Higgs sectors localized on matter curves of intersecting D7 branes. In [22], such scenarios were proposed as candidates for shift symmetric models. The crucial question - what is the moduli dependence of the Kähler metric pertaining to the $H_u H_d + c.c.$ terms - can not be answered in a straightforward fashion by matching to the dimensionally reduced action. In analogy to radion mediated SUSY breaking in 5D/6D models, we conjecture that the relevant part of the Kähler metric might be recovered by considering dimensionally reduced actions in the presence of warping. This may open up the possibility to explicitly fine-tune the Higgs mass through fluxes.

We have then analyzed scenarios with an exchange symmetry realized with D6 branes at angles [22]. A crucial feature of these models is that the diagonal entries in the Higgs mass matrix are generated solely via a supersymmetric mass term corresponding to a small separation of the brane stacks supporting the Higgs hypermultiplets. Due to the absence of soft Higgs masses at leading order, the scenario is automatically exchange symmetric. The challenge is then to generate a $B\mu$ term of equal size without switching on comparable diagonal soft masses. In models [56, 61] where D6 brane configurations break SUSY completely (a situation which is dual to SUSY breaking open string fluxes in IIB) such terms can be generated. However, the interpretation of a soft breaking parameter as an "F-term like" $B\mu$ term rather than diagonal "D-term like" soft masses rests entirely on identifying one particular 4D $\mathcal{N} = 1$ supersymmetry as the surviving one. This scheme only yields $\lambda = 0$ in the low energy theory if the Higgs quartic potential is precisely an MSSM-like D-term with respect to this surviving generator. We have shown how this can in principle be achieved in the simplified setting of D6 branes at angles in flat space. The key is that certain types of brane angles (non-factorizable) can actually correspond to F-terms.

Shift/exchange symmetry in the Higgs sector is necessarily broken at 1-loop level due to the Yukawa couplings to matter, in particular to the top quark. We have already given the resulting corrections to $\tan \beta = 1$ in [16] and found that they are small enough to maintain predictivity of an approximate shift/exchange symmetry. In this paper, we have also considered the loop corrections to the quartic coupling itself. The decoupling

and threshold contributions which arise when the heavy MSSM states are integrated out generically are of similar size as the shift violating corrections. Both types of corrections merely yield shifts of the physical Higgs mass of $|\Delta m_h| \lesssim 2$ GeV, often less. This is in contrast to TeV scale SUSY and is mainly due to the relative smallness of the top Yukawa coupling at high renormalization scales [15]. Furthermore, in scenarios with a small splitting between the soft parameters, some of the contributions tend to cancel. In summary, we find that in all but the extreme cases, the overall loop corrections to the quartic coupling are small, and the preferred region for the soft breaking scale remains close to the naive estimate.

Finally, we have considered whether the incomplete decoupling of F-term potentials (which we have tried to avoid until now) might in some situations provide us with a UV completion with $\lambda < 0$. We have shown, based on a simple NMSSM-like model, how such an unusual situation can arise. We found that the soft scale at which our UV completion comes in can be raised all the way up near the Planck scale even if the top quark mass turns out to be towards the heavy end of the currently favored range. While one naively expects that some new physics must come into play at μ_λ to avoid the instability of the scalar potential, this is not strictly speaking necessary: There may simply be an energy gap between μ_λ and m_{UV} (e.g. the string scale) where no stable 4d effective theory exists. Also, arguing for new physics at μ_λ from a cosmological point of view is not fully convincing: In the early history of the universe, even if it was very hot, the finite-temperature effective potential will presumably be stable independently of the running of λ. In the late history, all that matters is a sufficiently long lifetime of our (possibly metastable) minimum [9]. In order to match the SM running of the quartic coupling to such an unstable UV completion, we have addressed a different but related question, namely whether perturbation theory in the $\lambda < 0$ regime is stable enough such that λ is still a meaningful UV boundary conditions of the SM RGEs. We have given some simple arguments which suggest that the lifetime of "vacuum" states is sufficiently long to allow such a matching at least at the level of precision required for our purposes.

In this paper, we have not discussed dark matter nor flavor physics in any detail. As we have already noted in [16], and as was worked out in [22], the typical scales which appear in these high-scale SUSY scenarios are in an interesting range for Axion CDM as well as near typical Seesaw scales. Indeed, while it does not seem impossible to extend our models to variations of split SUSY [6] with an electroweak WIMP candidate at the TeV scale (however, see [81] for some Caveats), axions are ubiquitous in string compactifications and certainly provide the most compelling dark matter candiates in this class of models. Further research in this direction would certainly be interesting. Other recent proposals relating axions to high-scale physics can for example be found in [20, 82]. An interesting alternative to the MSSM or NMSSM like Higgs sectors in high-scale SUSY are single-Higgs SUSY models (see [21, 83]). For models employing shift symmetries of the type discussed in this paper in the context of LARGE volume scenarios, see [84].

Acknowledgments

Appendix

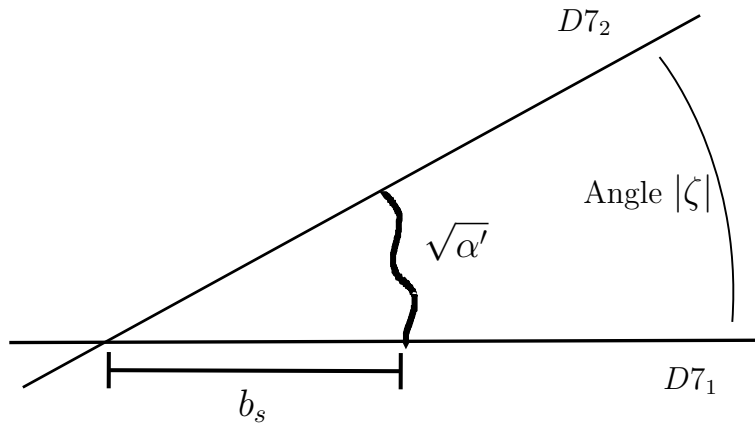## A   The transition from bulk matter to intersection matter

Figure 5: An illustration of the scales involved in the localization of string states in the bifundamental representation.

In section 2 we discuss the possibilities for shift symmetry in IIB models in which the Higgs sector is realized as D7 brane matter. We now want to see how matter on intersection curves arises as a deformation of this scenario in order to study the moduli dependence of the matter Kähler potential.

There is in principle a continuous transition from localized matter to matter on a brane stack if the compactification allows a brane deformation modulus ζ corresponding to a nontrivial profile which locally looks like $\phi \sim |\zeta| z_1$ of one of the adjoint brane scalars ϕ parameterizing the relative transverse brane movement (here, z_1 is one of the internal brane-parallel directions z_1, z_2, and the normalization of ζ is such that $\zeta \sim \mathcal{O}(1)$ corresponds to generic angles). However, the known moduli dependences of K on the dilaton S and Kähler moduli T_i differ between the two cases [32, 39]. It is thus our aim to understand at which point this transition between a "bulk-like" and a "matter-curve-like" situation occurs.

For simplicity we assume two spacetime filling D7 branes with parallel directions z_1 and z_2 in the extra dimensions. Our starting point is a situation where they are on top

of each other with the known bulk-matter Kähler metric [32, 39, 40, 42, 45]

$$K_{Bulk} \sim s^{-1} \tag{67}$$

and no dependence on t, where s and t are the real parts of the dilaton and relevant Kähler modulus respectively. We now switch on an adjoint vev

$$\langle \phi \rangle \sim \gamma z_1 \tag{68}$$

where the constant γ parameterizes the angle. The brane-matter coupled to this vev profile now also obtains a nontrivial profile

$$\psi(z_1, z_2) \sim e^{-\gamma |z_1|^2} f(z_2) \,. \tag{69}$$

Integrating over z_1 and z_2 in the kinetic term, we obtain

$$\int d^2 z_1 \int d^2 z_2 |\psi|^2 = \int d^2 z_2 |f|^2 \int d^2 z_1 e^{-\gamma |z_1|^2} \sim \frac{1}{\gamma} \tag{70}$$

Consequently, the kinetic part of the matter Kähler potential derived from the intersection curve scales like

$$K_{int} \sim s^{-1} \gamma^{-1} \,. \tag{71}$$

We thus have to determine the dependence of γ on the moduli. Let us first note that the typical extension of the "classical" particle wave functions in terms of complex geometry coordinates is

$$|z^{crit}| \sim \gamma^{-1/2} \,. \tag{72}$$

We now match this field theory quantity to the string picture illustrated in Figure 5. It is clear that the extension of the particle wave functions is given by the region of size b_s where the brane separation is small compared to the string length. In the string frame, the relation between z^{crit} and b_s is, up to proportionality factors, given by

$$b_s \sim |z_{crit}| R_s \tag{73}$$

Here, R_s is the typical radius of the compactification in the string frame which scales like the Kähler modulus. In the string frame, the string length is constant, and we can thus read off (see Fig. 5) that

$$|\zeta| b_s \sim \sqrt{\alpha'} \tag{74}$$

and consequently

$$\gamma^{-1/2} R_s \sim b_s \sim \frac{\sqrt{\alpha'}}{|\zeta|} \sim \frac{1}{|\zeta|} \,. \tag{75}$$

The relation between string frame and Einstein frame metric is given by

$$g_{ij}^s = g_{ij}^E g_s^{1/2} \tag{76}$$

and thus likewise for the Kähler modulus

$$R_s^2 \sim R_E^2 g_s^{1/2} \sim t^{1/2} g_s^{1/2} = t^{1/2} s^{-1/2} \,. \tag{77}$$

Putting everything together, we find that

$$K_{int} \sim \frac{1}{s}\frac{1}{\gamma} = \frac{1}{s}\frac{1}{R_s^2|\zeta|^2} = \frac{1}{\sqrt{st}}\frac{1}{|\zeta|^2} \tag{78}$$

This moduli dependence is consistent with what was found in [39]. We also obtain the dependence on the brane deformation modulus ζ.

When does the transition between bulk and intersection curve Kähler potential occur? We have assumed that the extent of the classical wave function is limited by the brane separation. In this situation, the wave function does not "feel" the overall volume of the cycle on which the brane is wrapped. This changes as soon as the extent of the classical wave function becomes comparable to the compactification radius, at $\gamma = 1$ or $|\zeta|^2 \sim \sqrt{s/t}$. This transition is captured in (23).

B   Reliability of effective field theory in the classically unstable regime of negative λ

In section 4.2 we discuss UV completions with negative quartic coupling and argue that it is possible to match them perturbatively to the SM at the soft breaking scale. Let us try to be a bit more precise: We consider a model with a massless scalar H_0 with negative quartic potential. This is not the full theory, but a good approximation as long as $|H_0| \ll M$. Furthermore, we pretend for the sake of the following qualitative analysis that H_0 is a single real scalar rather than a complex doublet.

To get rid of IR problems, we compactify the theory on a T^3 with volume R^3. The instability is associated only with the zero mode ϕ of H_0. Neglecting for the moment the quartic coupling, the dynamics of this mode corresponds simply to that of a free quantum mechanical particle with position $x = \phi R^2$ and mass $m = 1/R$,

$$S = \int dt\, R^3 \frac{1}{2}\dot{\phi}^2 = \int dt \frac{1}{2R}\dot{x}^2\,. \tag{79}$$

Let us take our initial state to be a Gaussian Schrödinger wave function with width δx_0 and momentum uncertainty $\sim 1/\delta x_0$. According to the Schrödinger equation, it will spread with time in the familiar way: $\delta x \sim \sqrt{\delta x_0^2 + (t-t_0)^2/4m^2\delta x_0^2}$. Returning to the field theory model, this translates into an unavoidable quantum uncertainty of the field position $\delta\phi$ and field velocity $\delta\dot{\phi} \sim 1/(R^3\delta\phi_0)$. The former then grows as

$$\delta\phi \sim \sqrt{\delta\phi_0^2 + (t-t_0)^2/4R^6\delta\phi_0^2}\,, \tag{80}$$

while the latter remains fixed.

We want to match this to the classical analysis of vacuum decay in the inverted quartic potential:

$$\rho = \frac{\dot{\phi}^2}{2} - |\lambda|\phi^4 \quad \Rightarrow \quad t = \int \frac{d\phi}{\sqrt{2(\rho + |\lambda|\phi^4)}}\,. \tag{81}$$

33

Here $\rho = E/R^3$ is the (conserved) energy density associated with the zero mode. One immediately observes two regimes: An early regime where

$$\phi = \phi_1 + (t - t_0)\sqrt{2\rho}\,, \tag{82}$$

and a late regime,

$$\phi = \frac{1}{\sqrt{2|\lambda|}(-t)} \qquad \text{with} \qquad t < 0\,, \tag{83}$$

diverging in finite time as expected. The matching is at the time or, equivalently, ϕ-value satisfying $|\lambda|\phi^4 \sim \rho$, with ρ a conserved quantity encoding the initial conditions.

We could now optimize our choice of quantum state (i.e. of $\delta\phi_0$) to live for as long as possible 'on the top of the hill'. Instead, we start with a simple-minded guess: Let $\delta\phi_0 \sim 1/R$, such that R remains the only dimensionful quantity in the problem, and attempt to match this immediately to the classical description. According to our Gaussian wave function[18] the typical classical field configuration has $\phi_1 \sim \delta\phi_0 \sim \pm 1/R$, velocity $\dot{\phi} \sim 1/R^2$ and energy density $\rho \sim 1/R^4$ (cf. (81), assuming also $|\lambda| \ll 1$). The field evolves classically (cf. (82) with $t_0 = 0$) as

$$\phi \sim \pm\frac{1}{R} + \frac{t}{R^2}\,, \tag{84}$$

reaching the dangerous late regime of fast decay at $\phi \sim 1/(|\lambda|^{1/4}R)$. The corresponding critical time is $t_c \sim R/|\lambda|^{1/4}$. Our early matching from quantum to classical is a posteriori justified by the observation that the same conclusion would have been reached in quantum evolution according to (80).

Thus, we can safely think in terms of a theory with IR cutoff $1/R$, asking any dynamical question which does not require time scales larger than t_c (with $t_c \gg R$ for $|\lambda| \ll 1$). The cutoff can be chosen as low as we wish. It will generically induce relevant operators such as an effective Higgs mass, but since these are suppressed relative to $1/R^2$ by a loop factor, they should not interfere with physics at time scales relevant for our argument. Thus, our scenario with negative high-scale λ works as long as λ runs to a positive value somewhere in the IR.

References

[1] G. Aad *et al.* [ATLAS Collaboration], "Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC," Phys. Lett. B **716**, 1 (2012) [arXiv:1207.7214 [hep-ex]].
For recent updates, see
The ATLAS collaboration, ATLAS-CONF-2013-014
The ATLAS collaboration, ATLAS-CONF-2013-034.

---

[18] Actually, we are of course dealing with a wave functional, but we can ignore the non-zero modes for our purposes.

[2] S. Chatrchyan *et al.* [CMS Collaboration], "Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC," Phys. Lett. B **716**, 30 (2012) [arXiv:1207.7235 [hep-ex]].
For recent updates (in particular of the $\gamma\gamma$ channel), see
The CMS collaboration, CMS PAS-HIG-12-045
The CMS collaboration, CMS PAS-HIG-13-001.

[3] M. Baak, M. Goebel, J. Haller, A. Hoecker, D. Ludwig, K. Moenig, M. Schott and J. Stelzer, "Updated Status of the Global Electroweak Fit and Constraints on New Physics," Eur. Phys. J. C **72** (2012) 2003 [arXiv:1107.0975 [hep-ph]].

[4] G. Isidori, Y. Nir and G. Perez, "Flavor Physics Constraints for Physics Beyond the Standard Model," Ann. Rev. Nucl. Part. Sci. **60** (2010) 355 [arXiv:1002.0900 [hep-ph]].

[5] F. Denef and M. R. Douglas, "Distributions of flux vacua," JHEP **0405** (2004) 072 [hep-th/0404116] and "Distributions of nonsupersymmetric flux vacua," JHEP **0503** (2005) 061 [hep-th/0411183].

[6] N. Arkani-Hamed and S. Dimopoulos, "Supersymmetric unification without low energy supersymmetry and signatures for fine-tuning at the LHC," JHEP **0506** (2005) 073 [hep-th/0405159];
G. F. Giudice and A. Romanino, "Split supersymmetry," Nucl. Phys. B **699** (2004) 65 [Erratum-ibid. B **706** (2005) 65] [hep-ph/0406088];
L. J. Hall and Y. Nomura, "Spread Supersymmetry," JHEP **1201** (2012) 082 [arXiv:1111.4519 [hep-ph]];
N. Arkani-Hamed, A. Gupta, D. E. Kaplan, N. Weiner and T. Zorawski, "Simply Unnatural Supersymmetry," arXiv:1212.6971 [hep-ph].

[7] M. Lindner, M. Sher and H. W. Zaglauer, "Probing Vacuum Stability Bounds at the Fermilab Collider," Phys. Lett. B **228** (1989) 139;
M. Sher, "Electroweak Higgs Potentials and Vacuum Stability," Phys. Rept. **179** (1989) 273;
J. A. Casas, J. R. Espinosa and M. Quiros, "Improved Higgs mass stability bound in the Standard Model and implications f or supersymmetry," Phys. Lett. B **342** (1995) 171 [hep-ph/9409458] and "Standard Model stability bounds for new physics within LHC reach," Phys. Lett. B **382** (1996) 374 [hep-ph/9603227];
C. D. Froggatt and H. B. Nielsen, "Standard Model criticality prediction: Top mass 173 +- 5-GeV and Higgs mass 135 +- 9-GeV," Phys. Lett. B **368** (1996) 96 [hep-ph/9511371];
C. D. Froggatt, H. B. Nielsen and Y. Takanishi, "Standard Model Higgs boson mass from borderline metastability of the vacuum," Phys. Rev. D **64** (2001) 113014 [hep-ph/0104161].

[8] M. E. Cabrera, J. A. Casas and A. Delgado, "Upper Bounds on Superpartner Masses from Upper Bounds on the Higgs Boson Mass," Phys. Rev. Lett. **108**, 021802 (2012) [arXiv:1108.3867 [hep-ph]].

[9] G. Degrassi, S. Di Vita, J. Elias-Miro, J. R. Espinosa, G. F. Giudice, G. Isidori and A. Strumia, "Higgs mass and vacuum stability in the Standard Model at NNLO," JHEP **1208** (2012) 098 [arXiv:1205.6497 [hep-ph]].

[10] M. Holthausen, K. S. Lim and M. Lindner, "Planck scale Boundary Conditions and the Higgs Mass," arXiv:1112.2415 [hep-ph];
C. Wetterich, "Where to look for solving the gauge hierarchy problem?," arXiv:1112.2910 [hep-ph].

[11] G. F. Giudice and A. Strumia, "Probing High-Scale and Split Supersymmetry with Higgs Mass Measurements," Nucl. Phys. B **858** (2012) 63 [arXiv:1108.6077v2 [hep-ph]].

[12] J. Elias-Miro, J. R. Espinosa, G. F. Giudice, G. Isidori, A. Riotto and A. Strumia, "Higgs mass implications on the stability of the electroweak vacuum," arXiv:1112.3022 [hep-ph].

[13] K. G. Chetyrkin and M. F. Zoller, "β-function for the Higgs self-interaction in the Standard Model at three-loop level," arXiv:1303.2890 [hep-ph].
M. F. Zoller, "Vacuum stability in the SM and the three-loop β-function for the Higgs self-interaction," arXiv:1209.5609 [hep-ph].
L. N. Mihaila, J. Salomon and M. Steinhauser, "Renormalization constants and beta functions for the gauge couplings of the Standard Model to three-loop order," Phys. Rev. D **86**, 096008 (2012) [arXiv:1208.3357 [hep-ph]].
K. G. Chetyrkin and M. F. Zoller, "Three-loop β-functions for top-Yukawa and the Higgs self-interaction in the Standard Model," JHEP **1206**, 033 (2012) [arXiv:1205.2892 [hep-ph]].
L. N. Mihaila, J. Salomon and M. Steinhauser, "Gauge Coupling Beta Functions in the Standard Model to Three Loops," Phys. Rev. Lett. **108**, 151602 (2012) [arXiv:1201.5868 [hep-ph]].

[14] For a review see e.g. S. P. Martin, "A Supersymmetry primer," In *Kane, G.L. (ed.): Perspectives on supersymmetry II* 1-153 [hep-ph/9709356].

[15] L. J. Hall and Y. Nomura, "A Finely-Predicted Higgs Boson Mass from A Finely-Tuned Weak Scale," JHEP **1003** (2010) 076 [arXiv:0910.2235 [hep-ph]];

[16] A. Hebecker, A. K. Knochel and T. Weigand, "A Shift Symmetry in the Higgs Sector: Experimental Hints and Stringy Realizations," JHEP **1206** (2012) 093 [arXiv:1204.2551 [hep-th]].

[17] G. Lopes Cardoso, D. Lüst and T. Mohaupt, "Moduli spaces and target space duality symmetries in (0,2) Z(N) orbifold theories with continuous Wilson lines," Nucl. Phys. B **432** (1994) 68 [arXiv:hep-th/9405002];
I. Antoniadis, E. Gava, K. S. Narain and T. R. Taylor, "Effective mu term in superstring theory," Nucl. Phys. B **432** (1994) 187 [arXiv:hep-th/9405024];
A. Brignole, L. E. Ibáñez, C. Muñoz and C. Scheich, "Some Issues In Soft Susy Breaking Terms From Dilaton / Moduli Sectors," Z. Phys. C **74** (1997) 157

[arXiv:hep-ph/9508258];
A. Brignole, L. E. Ibáñez and C. Muñoz, "Orbifold-induced mu term and electroweak symmetry breaking," Phys. Lett. B **387** (1996) 769 [arXiv:hep-ph/9607405];

[18] A. Brignole, L. E. Ibáñez and C. Muñoz, "Soft supersymmetry-breaking terms from supergravity and superstring models," arXiv:hep-ph/9707209. L. J. Dixon, V. Kaplunovsky and J. Louis, "On Effective Field Theories Describing (2,2) Vacua of the Heterotic String," Nucl. Phys. B **329**, 27 (1990).

[19] I. Gogoladze, N. Okada and Q. Shafi, "Higgs boson mass from gauge-Higgs unification," Phys. Lett. B **655** (2007) 257 [arXiv:0705.3035 [hep-ph]] and "Window For Higgs Boson Mass From Gauge-Higgs Unification," Phys. Lett. B **659** (2008) 316 [arXiv:0708.2503 [hep-ph]].

[20] M. Redi and A. Strumia, "Axion-Higgs Unification," JHEP **1211** (2012) 103 [arXiv:1208.6013 [hep-ph]].

[21] J. Unwin, "R-symmetric High Scale Supersymmetry," Phys. Rev. D **86** (2012) 095002 [arXiv:1210.4936 [hep-ph]].

[22] L. E. Ibáñez, F. Marchesano, D. Regalado and I. Valenzuela, "The Intermediate Scale MSSM, the Higgs Mass and F-theory Unification," JHEP **1207** (2012) 195 [arXiv:1206.2655 [hep-ph]].

[23] L. E. Ibáñez and I. Valenzuela, "The Higgs Mass as a Signature of Heavy SUSY," arXiv:1301.5167 [hep-ph].

[24] J. Elias-Miro, J. R. Espinosa, G. F. Giudice, H. M. Lee and A. Strumia, "Stabilization of the Electroweak Vacuum by a Scalar Threshold Effect," JHEP **1206**, 031 (2012) [arXiv:1203.0237 [hep-ph]].
L. A. Anchordoqui, I. Antoniadis, H. Goldberg, X. Huang, D. Lüst, T. R. Taylor and B. Vlcek, "Vacuum Stability of Standard Model^{++}," JHEP **1302**, 074 (2013) [arXiv:1208.2821 [hep-ph]].

[25] M. Cicoli, S. de Alwis and A. Westphal, "Heterotic Moduli Stabilization," arXiv:1304.1809 [hep-th].

[26] W. Buchmüller, K. Hamaguchi, O. Lebedev and M. Ratz, "Supersymmetric Standard Model from the heterotic string," Phys. Rev. Lett. **96**, 121602 (2006) [arXiv:hep-ph/0511035].
O. Lebedev, H. P. Nilles, S. Raby, S. Ramos-Sanchez, M. Ratz, P. K. S. Vaudrevange and A. Wingerter, "A mini-landscape of exact MSSM spectra in heterotic orbifolds," Phys. Lett. B **645** (2007) 88 [arXiv:hep-th/0611095].

[27] K. w. Choi, N. y. Haba, K. S. Jeong, K. i. Okumura, Y. Shimizu and M. Yamaguchi, "Electroweak symmetry breaking in supersymmetric gauge-Higgs unification models," JHEP **0402** (2004) 037 [arXiv:hep-ph/0312178];
A. Hebecker, J. March-Russell and R. Ziegler, "Inducing the mu and the B mu

Term by the Radion and the 5d Chern-Simons Term," JHEP **0908** (2009) 064 [arXiv:0801.4101 [hep-ph]];
F. Brümmer, S. Fichet, A. Hebecker and S. Kraml, "Phenomenology of Supersymmetric Gauge-Higgs Unification," JHEP **0908** (2009) 011 [arXiv:0906.2957 [hep-ph]].

[28] F. Brümmer, R. Kappl, M. Ratz and K. Schmidt-Hoberg, "Approximate R-symmetries and the mu term," JHEP **1004** (2010) 006 [arXiv:1003.0084 [hep-th]];
F. Brümmer, S. Fichet, S. Kraml and R. K. Singh, "On SUSY GUTs with a degenerate Higgs mass matrix," JHEP **1008**, 096 (2010) [arXiv:1007.0321 [hep-ph]].

[29] D. K. M. Pena, H. P. Nilles and P. -K. Oehlmann, "A Zip-code for Quarks, Leptons and Higgs Bosons," JHEP **1212**, 024 (2012) [arXiv:1209.6041 [hep-th]].

[30] L. E. Ibáñez, H. P. Nilles and F. Quevedo, "Reducing The Rank Of The Gauge Group In Orbifold Compactifications Of The Heterotic String," Phys. Lett. B **192**, 332 (1987).

[31] G. Aldazabal, A. Font, L. E. Ibáñez and G. Violero, "D = 4, N=1, type IIB orientifolds," Nucl. Phys. B **536** (1998) 29 [hep-th/9804026].
R. Blumenhagen, V. Braun, T. W. Grimm and T. Weigand, "GUTs in Type IIB Orientifold Compactifications," Nucl. Phys. B **815** (2009) 1 [arXiv:0811.2936 [hep-th]].

[32] H. Jockers and J. Louis, "The Effective action of D7-branes in N = 1 Calabi-Yau orientifolds," Nucl. Phys. B **705** (2005) 167 [hep-th/0409098].

[33] A. Hebecker, S. C. Kraus, M. Kuntzler, D. Lüst and T. Weigand, "Fluxbranes: Moduli Stabilisation and Inflation," JHEP **1301** (2013) 095 [arXiv:1207.2766 [hep-th]].

[34] M. Arends, A. Hebecker, K. Heimpel, S. C. Kraus, D. Lüst, C. Mayrhofer, C. Schick, and T. Weigand, "Flat Directions in D7-Brane Moduli Space and Inflationary Model Building", work in progress.

[35] M. Kerstan and T. Weigand, "The Effective action of D6-branes in N=1 type IIA orientifolds," JHEP **1106** (2011) 105 [arXiv:1104.2329 [hep-th]].
T. W. Grimm and D. V. Lopes, "The N=1 effective actions of D-branes in Type IIA and IIB orientifolds," Nucl. Phys. B **855** (2012) 639 [arXiv:1104.2328 [hep-th]].

[36] R. Bohm, H. Gunther, C. Herrmann and J. Louis, "Compactification of type IIB string theory on Calabi-Yau threefolds," Nucl. Phys. B **569** (2000) 229 [hep-th/9908007];
T. W. Grimm and J. Louis, "The Effective action of type IIA Calabi-Yau orientifolds," Nucl. Phys. B **718** (2005) 153 [hep-th/0412277];
R. D'Auria, S. Ferrara and M. Trigiante, "C - map, very special quaternionic geometry and dual Kahler spaces," Phys. Lett. B **587** (2004) 138 [hep-th/0401161];
S. Alexandrov, B. Pioline, F. Saueressig and S. Vandoren, "D-instantons and twistors," JHEP **0903** (2009) 044 [arXiv:0812.4219 [hep-th]].

[37] R. Donagi and M. Wijnholt, "Gluing Branes II: Flavour Physics and String Duality," arXiv:1112.4854 [hep-th].

[38] A. Maharana and E. Palti, "Models of Particle Physics from Type IIB String Theory and F-theory: A Review," arXiv:1212.0555 [hep-th].
T. Weigand, "Lectures on F-theory compactifications and model building," Class. Quant. Grav. **27**, 214004 (2010) [arXiv:1009.3497 [hep-th]].

[39] L. Aparicio, D. G. Cerdeno and L. E. Ibáñez, "Modulus-dominated SUSY-breaking soft terms in F-theory and their test at LHC," JHEP **0807** (2008) 099 [arXiv:0805.2943 [hep-ph]].

[40] J. P. Conlon, D. Cremades and F. Quevedo, "Kahler potentials of chiral matter fields for Calabi-Yau string compactifications," JHEP **0701** (2007) 022 [hep-th/0609180].

[41] C. Beasley, J. J. Heckman and C. Vafa, "GUTs and Exceptional Branes in F-theory - I," JHEP **0901** (2009) 058 [arXiv:0802.3391 [hep-th]].

[42] A. Font and L. E. Ibáñez, "Matter wave functions and Yukawa couplings in F-theory Grand Unification," JHEP **0909** (2009) 036 [arXiv:0907.4895 [hep-th]];
J. P. Conlon and E. Palti, "Aspects of Flavour and Supersymmetry in F-theory GUTs," JHEP **1001** (2010) 029 [arXiv:0910.2413 [hep-th]];
F. Marchesano, P. McGuirk and G. Shiu, "Chiral matter wavefunctions in warped compactifications," JHEP **1105** (2011) 090 [arXiv:1012.2759 [hep-th]].

[43] L. E. Ibáñez, C. Munoz and S. Rigolin, "Aspect of type I string phenomenology," Nucl. Phys. B **553** (1999) 43 [hep-ph/9812397].

[44] D. Lüst, P. Mayr, R. Richter and S. Stieberger, "Scattering of gauge, matter, and moduli fields from intersecting branes," Nucl. Phys. B **696** (2004) 205 [hep-th/0404134];
D. Lüst, S. Reffert and S. Stieberger, "Flux-induced soft supersymmetry breaking in chiral type IIB orientifolds with D3 / D7-branes," Nucl. Phys. B **706** (2005) 3 [hep-th/0406092].

[45] T. Kawano, Y. Tsuchiya and T. Watari, "A Note on Kahler Potential of Charged Matter in F-theory," Phys. Lett. B **709**, 254 (2012) [arXiv:1112.2987 [hep-th]].

[46] Z. Chacko and M. A. Luty, "Radion mediated supersymmetry breaking," JHEP **0105** (2001) 067 [hep-ph/0008103];
D. E. Kaplan and N. Weiner, "Radion mediated supersymmetry breaking as a Scherk-Schwarz theory," hep-ph/0108001.

[47] D. Marti and A. Pomarol, "Supersymmetric theories with compact extra dimensions in N=1 superfields," Phys. Rev. D **64** (2001) 105025 [hep-th/0106256].

[48] P. Fayet, "Fermi-Bose Hypersymmetry," Nucl. Phys. B **113** (1976) 135.

[49] J. A. Strathdee, "Extended Poincare Supersymmetry," Int. J. Mod. Phys. A **2** (1987) 273.

[50] A. Hebecker and A. Westphal, "Gauge unification in extra dimensions: Power corrections vs. higher-dimension operators," Nucl. Phys. B **701** (2004) 273 [hep-th/0407014].

[51] R. Blumenhagen, M. Cvetič and T. Weigand, "Spacetime instanton corrections in 4D string vacua: The Seesaw mechanism for D-Brane models," Nucl. Phys. B **771** (2007) 113 [hep-th/0609191].
L. E. Ibáñez and A. M. Uranga, "Neutrino Majorana Masses from String Theory Instanton Effects," JHEP **0703** (2007) 052 [hep-th/0609213].
M. Buican, D. Malyshev, D. R. Morrison, H. Verlinde and M. Wijnholt, "D-branes at Singularities, Compactification, and Hypercharge," JHEP **0701** (2007) 107 [hep-th/0610007].
L. E. Ibáñez and A. M. Uranga, "Instanton induced open string superpotentials and branes at singularities," JHEP **0802** (2008) 103 [arXiv:0711.1316 [hep-th]].
M. Berg, J. P. Conlon, D. Marsh and L. T. Witkowski, "Superpotential desequestering in string models," JHEP **1302** (2013) 018 [arXiv:1207.1103 [hep-th]].

[52] P. Breitenlohner and M. F. Sohnius, "Matter coupling and nonlinear σ models in N=2 supergravity," Nucl. Phys. B **187** (1981) 409.

[53] M. F. Sohnius, "Introducing Supersymmetry," Phys. Rept. **128** (1985) 39.

[54] S. B. Giddings, S. Kachru and J. Polchinski, "Hierarchies from fluxes in string compactifications," Phys. Rev. D **66** (2002) 106006 [hep-th/0105097].

[55] C. P. Burgess, P. G. Camara, S. P. de Alwis, S. B. Giddings, A. Maharana, F. Quevedo and K. Suruliz, "Warped Supersymmetry Breaking," JHEP **0804** (2008) 053 [hep-th/0610255].

[56] D. Cremades, L. E. Ibáñez and F. Marchesano, "Intersecting brane models of particle physics and the Higgs mechanism," JHEP **0207** (2002) 022 [hep-th/0203160].

[57] R. Blumenhagen, M. Cvetič, P. Langacker and G. Shiu, "Toward realistic intersecting D-brane models," Ann. Rev. Nucl. Part. Sci. **55**, 71 (2005) [hep-th/0502005].
R. Blumenhagen, B. Kors, D. Lüst and S. Stieberger, "Four-dimensional String Compactifications with D-Branes, Orientifolds and Fluxes," Phys. Rept. **445** (2007) 1 [hep-th/0610327].
L. E. Ibáñez and A. M. Uranga, "String theory and particle physics: An introduction to string phenomenology," Cambridge, UK: Univ. Pr. (2012) 673 p

[58] M. Cvetič, G. Shiu and A. M. Uranga, "Chiral four-dimensional N=1 supersymmetric type 2A orientifolds from intersecting D6 branes," Nucl. Phys. B **615** (2001) 3 [hep-th/0107166].

[59] G. Honecker and T. Ott, "Getting just the supersymmetric standard model at intersecting branes on the Z(6) orientifold," Phys. Rev. D **70**, 126010 (2004) [Erratum-ibid. D **71**, 069902 (2005)] [hep-th/0404055].

[60] R. Blumenhagen, L. Goerlich, B. Kors and D. Lüst, "Noncommutative compactifications of type I strings on tori with magnetic background flux," JHEP **0010** (2000) 006 [hep-th/0007024].
G. Aldazabal, S. Franco, L. E. Ibáñez, R. Rabadan and A. M. Uranga, "Intersecting brane worlds," JHEP **0102** (2001) 047 [hep-ph/0011132].

[61] L. E. Ibáñez, F. Marchesano and R. Rabadan, "Getting just the standard model at intersecting branes," JHEP **0111** (2001) 002 [hep-th/0105155].

[62] R. Blumenhagen, B. Kors, D. Lüst and T. Ott, "The standard model from stable intersecting brane world orbifolds," Nucl. Phys. B **616** (2001) 3 [hep-th/0107138].

[63] R. Kallosh, "N=2 supersymmetry and de Sitter space," hep-th/0109168;
R. Kallosh and A. D. Linde, "P term, D term and F term inflation," JCAP **0310** (2003) 008 [hep-th/0306058].

[64] E. A. Mirabelli and M. E. Peskin, "Transmission of supersymmetry breaking from a four-dimensional boundary," Phys. Rev. D **58**, 065002 (1998) [hep-th/9712214].
C. Bouchart, A. Knochel and G. Moreau, "Discriminating 4D supersymmetry from its 5D warped version," Phys. Rev. D **84**, 015016 (2011) [arXiv:1101.0634 [hep-ph]].

[65] P. Anastasopoulos, I. Antoniadis, K. Benakli, M. D. Goodsell and A. Vichi, "One-loop adjoint masses for non-supersymmetric intersecting branes," JHEP **1108** (2011) 120 [arXiv:1105.0591 [hep-th]].

[66] M. Berkooz, M. R. Douglas and R. G. Leigh, "Branes intersecting at angles," Nucl. Phys. B **480** (1996) 265 [hep-th/9606139].

[67] Y. Okada, M. Yamaguchi and T. Yanagida, "Upper bound of the lightest Higgs boson mass in the minimal supersymmetric standard model," Prog. Theor. Phys. **85**, 1 (1991).
Y. Okada, M. Yamaguchi and T. Yanagida, "Renormalization group analysis on the Higgs mass in the softly broken supersymmetric standard model," Phys. Lett. B **262**, 54 (1991).

[68] A. Kusenko, V. Kuzmin and I. I. Tkachev, Phys. Lett. B **432**, 361 (1998) [hep-ph/9801405].
J. M. Cornwall, A. Kusenko, L. Pearce and R. D. Peccei, arXiv:1210.6433 [hep-ph].

[69] J. Beringer et al. (Particle Data Group), PR D86, 010001 (2012)
(URL: http://pdg.lbl.gov)

[70] A. Dobado, M. J. Herrero, W. Hollik and S. Penaranda, "Selfinteractions of the lightest MSSM Higgs boson in the large pseudoscalar mass limit," Phys. Rev. D **66**, 095016 (2002) [hep-ph/0208014].

[71] N. Arkani-Hamed, T. Gregoire and J. G. Wacker, "Higher dimensional supersymmetry in 4-D superspace," JHEP **0203**, 055 (2002) [hep-th/0101233]. A. Hebecker, "5-D superYang-Mills theory in 4-D superspace, superfield brane operators, and applications to orbifold GUTs," Nucl. Phys. B **632**, 101 (2002) [hep-ph/0112230].

[72] E. Halyo, hep-th/0405269.
R. Kallosh and A. D. Linde, "P term, D term and F term inflation," JCAP **0310**, 008 (2003) [hep-th/0306058].

[73] A. H. Guth and S. Y. Pi, "Fluctuations in the New Inflationary Universe," Phys. Rev. Lett. **49**, 1110 (1982);
A. A. Starobinsky, "Dynamics of Phase Transition in the New Inflationary Universe Scenario and Generation of Perturbations," Phys. Lett. B **117** (1982) 175;
A. Vilenkin, "Topological inflation," Phys. Rev. Lett. **72** (1994) 3137 [hep-th/9402085];
L. Boubekeur and D. .H. Lyth, "Hilltop inflation," JCAP **0507** (2005) 010 [hep-ph/0502047].

[74] A. Vilenkin, L. H. Ford "Gravitational Effects upon Cosmological Phase Transitions," Phys. Rev. D **26** (1982) 1231.

[75] A. A. Starobinsky, "Stochastic De Sitter (inflationary) Stage In The Early Universe," In *De Vega, H.j. (Ed.), Sanchez, N. (Ed.): Field Theory, Quantum Gravity and Strings*, 107-126;
M. Sasaki, H. Suzuki, K. Yamamoto, J. 'i. Yokoyama, "Superexpansionary divergence: Breakdown of perturbative quantum field theory in space-time with accelerated expansion," Class. Quant. Grav. **10** (1993) L55;
A. A. Starobinsky, J. Yokoyama, "Equilibrium state of a selfinteracting scalar field in the De Sitter background," Phys. Rev. D **50** (1994) 6357 [astro-ph/9407016].

[76] S. A. Abel, C. -S. Chu, J. Jaeckel and V. V. Khoze, "SUSY breaking by a metastable ground state: Why the early universe preferred the non-supersymmetric vacuum," JHEP **0701** (2007) 089 [hep-th/0610334].

[77] T. Plehn and M. Rauch, "Higgs Couplings after the Discovery," Europhys. Lett. **100**, 11002 (2012) [arXiv:1207.6108 [hep-ph]].
A. Djouadi and G. g. Moreau, "The couplings of the Higgs boson and its CP properties from fits of the signal strengths and their ratios at the 7+8 TeV LHC," arXiv:1303.6591 [hep-ph].

[78] M. Shaposhnikov and C. Wetterich, "Asymptotic safety of gravity and the Higgs boson mass," Phys. Lett. B **683** (2010) 196 [arXiv:0912.0208 [hep-th]].

[79] P. Nath, "Higgs Physics and Supersymmetry," Int. J. Mod. Phys. A **27** (2012) 1230029 [arXiv:1210.0520 [hep-ph]] and arXiv:1302.1863 [hep-ph].
P. Nath, "Perspectives on Higgs Boson and Supersymmetry," arXiv:1302.1863 [hep-ph].

[80] A. Djouadi and J. Quevillon, "The MSSM Higgs sector at a high M_{SUSY}: reopening the low $\tan\beta$ regime and the search for heavy Higgsses," arXiv:1304.1787 [hep-ph].

[81] A. Arvanitaki, N. Craig, S. Dimopoulos and G. Villadoro, "Mini-Split," arXiv:1210.0555 [hep-ph];

[82] A. Chatzistavrakidis, E. Erfani, H. P. Nilles and I. Zavala, "Axiology," JCAP **1209**, 006 (2012) [arXiv:1207.1128 [hep-ph]].

[83] R. Davies, J. March-Russell and M. McCullough, "A Supersymmetric One Higgs Doublet Model," JHEP **1104** (2011) 108 [arXiv:1103.1647 [hep-ph]].

[84] T. Higaki and F. Takahashi, "Dark Radiation and Dark Matter in Large Volume Compactifications," JHEP **1211**, 125 (2012) [arXiv:1208.3563 [hep-ph]].
M. Cicoli, J. P. Conlon and F. Quevedo, "Dark Radiation in LARGE Volume Models," Phys. Rev. D **87**, 043520 (2013) [arXiv:1208.3562 [hep-ph]].
T. Higaki, K. Kamada and F. Takahashi, "Higgs, Moduli Problem, Baryogenesis and Large Volume Compactifications," JHEP **1209**, 043 (2012) [arXiv:1207.2771 [hep-ph]].